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Flipped SU(5): a Grand Unified Superstring Theory (GUST) Prototype

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Abstract

In this Lecture we first review the basic properties that make flipped SU(5) a very economical and well motivated unified model, including some recent developments regarding the prediction for α_s and new implications for proton decay. Then we sketch the derivation of flipped SU(5) from strings, stressing some new results concerning the cosmological constant, the stability of the no-scale mechanism, and a new mechanism for generating the "LEP" scale $M_{\rm LEP} \sim 10^{16}$ GeV in string models. Finally we present a sample supersymmetry breaking scenario, where all sparticle masses depend on a single parameter, and discuss the experimental signatures at LEP and the Tevatron. This scenario also predicts the top-quark mass to be close to 175 GeV.

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1 Standard Flipped SU(5)

1.1 Status of GUTs

In an nutshell, the status of GUTs can be summarized as follows. LEP has measured the Standard Model gauge couplings with unprecedented precision, leading to the following world averages

$$\sin^2\theta_W = 0.23143 \pm 0.00028 \, , \tag{1}$$
$$\alpha_s(M_Z) = 0.118 \pm 0.006 \, . \tag{2}$$

These couplings run with energy scale as prescribed by the renormalization group equations (RGEs), and may or may not converge at a single point. In the simplest GUTs without supersymmetry (*e.g.*, the Georgi-Glashow SU(5) model [1]) the gauge couplings do not converge [2] and thus fail the defining property of GUTs; a result that confirmed their earlier demise due to their incorrect proton lifetime prediction. Moreover, simple or complicated GUTs without supersymmetry are undesirable, as they fail to resolve the gauge hierarchy problem, or break dynamically the electroweak symmetry via radiative corrections. On the other hand, in the simplest SUSY GUTs the gauge couplings do converge around $M_{\rm GUT} \sim 10^{16}\,{\rm GeV}$ [3]. This contrast is illustrated in Fig. 1. The question then becomes

Are all GUTs with supersymmetry equally viable?

The answer to this question has to be no, as can be illustrated in the case of the minimal SU(5) SUSY GUT, which has several shortcomings. Some of these tend to be overlooked, as the problem with the doublet-triplet splitting, the lack of neutrino masses, and the lack of a mechanism for baryogenesis. More recently it has become apparent that its prediction for $\alpha_s(M_Z)$ [4]

$$\mathrm{SU}(5): \qquad \alpha_s(M_Z) > 0.123 \tag{3}$$

is very problematic. These troubles are tabulated below, and contrasted with their counterparts in Flipped SU(5) [5, 6]. Some people like to perform "model indepen-

Table 1: Comparison between SU(5) and flipped SU(5) GUT features.

Basic GUT tests	SU(5)	Flipped SU(5)
$\sin^2\theta_W \Rightarrow \alpha_3(M_Z)$	$\times$	$\surd$
Proton decay	$p \to \bar\nu K^+$	$p \to e^+\pi^0$
Doublet-triplet splitting	$\times$	$\surd$
Neutrino masses	$\times$	$\surd$
Baryogenesis	$\times$	$\surd$

dent" analyses by ignoring the GUT structure altogether, as is the case in "MSSM"

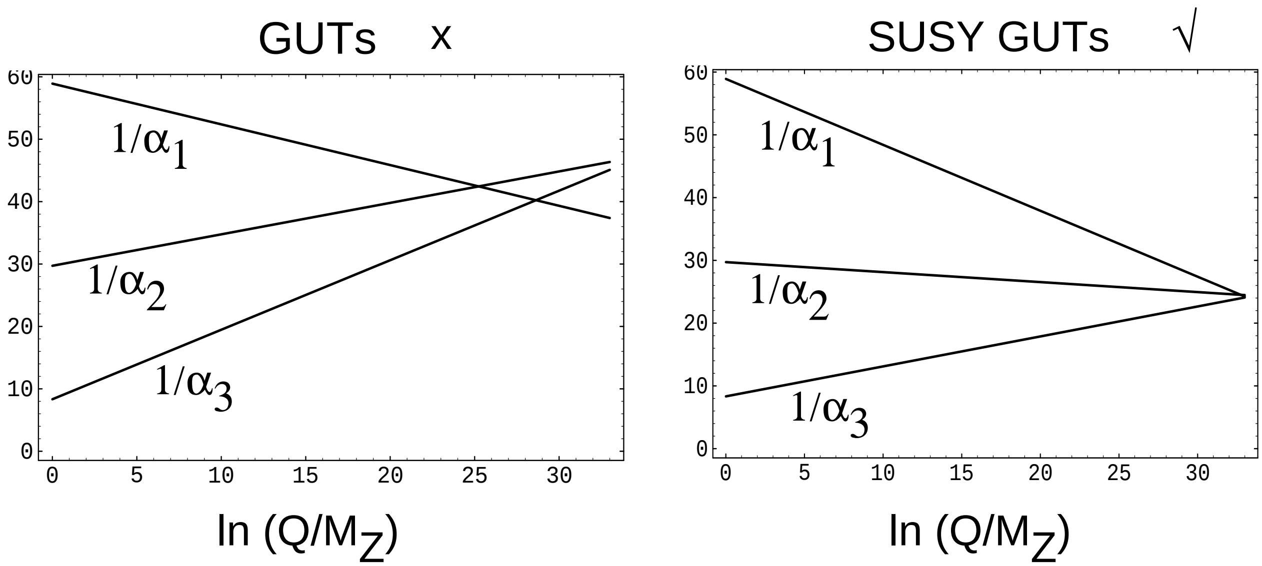

Figure 1: Running of the Standard Model gauge couplings in minimal GUTs and SUSY GUTs, showing the failure of the former and the success of the latter.

studies. This makes little sense, as the seed of the low-energy model predictions is found at or above the GUT scale. For instance, the constraint from proton decay is many times neglected, as if it could be satisfied automatically. Yet, in GUT models where this does happen (as in flipped SU(5) [6]) the GUT structure is quite different from that assumed in these "generic" models. When considering different SUSY GUT models, one can also ask the question

What really is the "minimal" GUT?

The analysis of Georgi and Glashow indicated that the smallest GUT structure containing the non-abelian Standard Model gauge groups is SU(5). Judging by the number of group generators, one would then consider flipped SU(5), and afterwards SO(10). However, we would argue that SU(5) is out of the running, as it lacks the desirable features mentioned above. Flipped SU(5) is the next larger gauge structure, and thus is the "minimal" SUSY GUT incorporating all the desirable GUT features mentioned above.[1] Flipped SU(5) differs from SU(5) and SO(10) in a very important respect: the electric charge generator does not lie completely within SU(5), as the gauge group is really SU(5)$\times$U(1). We would like to argue that such "non-grand" unification of the non-abelian Standard Model gauge couplings is in fact desirable, as the complete "grand" unification of all interactions must also include the gravitational sector (*i.e.*, the hidden sector and the supersymmetry breaking sector), and this is predicted not to occur until the string scale. To make the relevance of this

[1]These features can be incorporated in the missing-doublet version (MDM) of SU(5) [7], but this model is quite cumbersome and hardly "minimal" [8].

Table 2: Comparison between unified and grand unified SU(2) and SU(5) gauge groups and their properties.

SU(2) [Georgi-Glashow '72] • "grand" unified • $W^{\pm}, \gamma$: γ inside SU(2) • Higgs triplet (adjoint) • Neutral currents exist (1973) • Wrong!	SU(2)×U(1) [Glashow '61, Weinberg '67, Salam '68] • unified • $W^{\pm}, Z, \gamma$: $\gamma = \{W^3$ [SU(2)], B [U(1)]$\}$ • Higgs doublet (antisymmetric) • SU(3) not accounted for; grand unification later • Right!
SU(5) [Georgi-Glashow '74] • grand unified • $W^{\pm}, W^3, B, X, Y$ γ inside SU(5) • Higgs **24** (adjoint) • $\alpha_3 > 0.123$ $M_{\rm GUT} \sim \frac{1}{100} M_{\rm Pl}$ • Wrong?	SU(5)×U(1) [Barr '82, Derendinger-Kim-Nanopoulos '84, Antoniadis-Ellis-Hagelin-Nanopoulos '87] • unified • $W^{\pm}, W^3, B, X, Y, \widetilde{B}$ γ: (W^3, B) [SU(5)], $\widetilde{B}$ [U(1)] • Higgs $\mathbf{10}$,$\overline{\mathbf{10}}$ (antisymmetric) • Gravity* not accounted for; grand unification later • Right?

* "Gravity" = supersymmetry breaking, hidden sector gauge groups, string unification.

argument more apparent, we have contrasted the cases of unified and grand unified SU(2) and SU(5) gauge groups in Table 2.

1.2 Basic flipped SU(5) features

Because of the matter field content and gauge structure in flipped SU(5), several desirable GUT mechanisms are present [6] and work rather naturally, as we know briefly review.

1.2.1 Matter fields

The model includes three generations of quarks and leptons; ($\mathbf{10}$,$\overline{\mathbf{10}}$) Higgs GUT representations whose $\nu^c_H, \nu^c_{\bar H}$ components break the SU(5)×U(1) GUT symmetry down to the Standard Model once they acquire suitable vevs; and a pair of Higgs pentaplets which include the two light Higgs doublets. Some singlet fields (ϕ) are also present.

$$F_{(10)} = \{Q, d^c, \nu^c\};\ \bar f_{(\bar 5)} = \{L, u^c\};\ l_{(1)} = e^c \quad \text{(3 generations)}$$

$$H_{(10)} = \{Q_H, d^c_H, \nu^c_H\};\ \bar H_{(\overline{10})} = \{Q_{\bar H}, d^c_{\bar H}, \nu^c_{\bar H}\}$$

$$h_{(5)} = \{H_2, H_3\};\ \bar h_{(\bar 5)} = \{\bar H_2, \bar H_3\}$$

1.2.2 GUT superpotential

As allowed by the SU(5)×U(1) gauge symmetry, GUT couplings include interactions to effect the missing-partner mechanism and the see-saw mechanism

$$W_G = \lambda_4 HHh + \lambda_5 \bar{H}\bar{H}\bar{h} + \lambda_6 F\bar{H}\phi + \mu h\bar{h} \tag{4}$$

1.2.3 Doublet-triplet splitting

The components of the Higgs pentaplets must be split

$$h = \left(\begin{array}{c} H_2 \\ H_3 \end{array} \right) \begin{array}{l} \text{electroweak symmetry breaking} \\ \text{proton decay} \end{array} \tag{5}$$

because of their very different roles. The interactions in W_G

$$\lambda_4 HHh \rightarrow \lambda_4 d^c_H \langle \nu^c_H \rangle H_3 \tag{6}$$

$$\lambda_5 \bar{H}\bar{H}\bar{h} \rightarrow \lambda_5 \bar{d}^c_H \langle \bar{\nu}^c_H \rangle \bar{H}_3 \tag{7}$$

make the triplets heavy, while leaving the doublets light (the missing partner mechanism). A similar mechanism in the missing-doublet extension of SU(5) (MDM) requires the introduction of large representations ($\mathbf{50}$,$\overline{\mathbf{50}}$,$\mathbf{75}$) for this sole purpose [7].

1.2.4 Yukawa superpotential

The "flipping" of the assignments of the Standard Model fields to the SU(5) representations entails the couplings

$$\lambda_u F\bar{f}\,\bar{h} + \lambda_d FFh + \lambda_e \bar{f} l^c h\ . \tag{8}$$

Note that the usual $\lambda_b = \lambda_\tau$ relation in SU(5) is replaced by $\lambda_t = \lambda_\nu$, which has consequences for the spectrum of neutrino masses.

1.2.5 Neutrino masses

The "flipping" mentioned above brings the ν^c field into the F representation, providing a source of Dirac neutrino masses, in addition to the see-saw type coupling in W_G. These interactions result in a 3×3 see-saw mass matrix with the high scale provided by $M_U = \lambda_6 \langle \nu^c_H \rangle$,

$$\left. \begin{array}{l} \lambda_u F\bar{f}h \rightarrow m_u \nu\nu^c \\ \lambda_6 F\bar{H}\phi \rightarrow \lambda_6 \langle \nu^c_{\bar{H}} \rangle \nu^c\,\phi \end{array} \right\} \qquad M_\nu = \begin{array}{c} \\ \nu \\ \nu^c \\ \phi \end{array} \begin{array}{c} \begin{array}{ccc} \nu & \nu^c & \phi \end{array} \\ \left(\begin{array}{ccc} 0 & m_u & 0 \\ m_u & 0 & M_U \\ 0 & M_U & M \end{array} \right) \end{array} \tag{9}$$

The light neutrino masses then become

$$m_{\nu_{e,\mu,\tau}} \sim \frac{m_{u,c,t}^2}{M_U^2/M} \tag{10}$$

With the typical values $M_U \sim 10^{15}\,\mathrm{GeV}$ and $M \sim 10^{18}\,\mathrm{GeV}$ we get $m_{\nu_\tau} \sim 10\,\mathrm{eV}$ [9], which provides a very desirable source of hot dark matter. The result for $m_{\nu_\mu} \sim 10^{-3}\,\mathrm{eV}$ is consistent with the MSW explanation for the solar neutrino deficit [9]. Also, the right-handed (or "flipped") neutrino provides a source of lepton asymmetry which is later recycled into a baryon asymmetry by the electroweak-scale sphaleron interactions [10].

1.2.6 Proton decay

Dimension-six operators mediate proton decay via the exchange of the X, Y GUT gauge bosons, leading to the dominant decay mode $p \to e^+\pi^0$. Whether this mode is observable at SuperKamiokande (starting in mid 1996) or not depends on the magnitude of the unification scale. Quantitative predictions are given in Section 1.3 below. Proton decay via dimension-five operators is typically the largest process in traditional GUTs, such as minimal SU(5), where experimental limits on the mode $p \to \bar\nu K^+$ impose strict restrictions on the allowed parameter space of the model. The elementary diagram responsible for this reaction is shown in Fig. 2. The vertices originate from pieces of the couplings FFh, $F\bar f\bar h$

$$\lambda_d FFh \supset QQH_3\,; \qquad \lambda_u F\bar f\,\bar h \supset QL\bar H_3 \tag{11}$$

and are therefore proportional to the Yukawa matrices. Once dressed by squark and gaugino loops, this diagram leads to the well known $p \to \bar\nu K^+$ decay mode. However, in flipped SU(5) this does not happen because there is no $H_3, \bar H_3$ mixing, even though $H_3, \bar H_3$ are individually heavy via the doublet-triplet splitting mechanism discussed above.

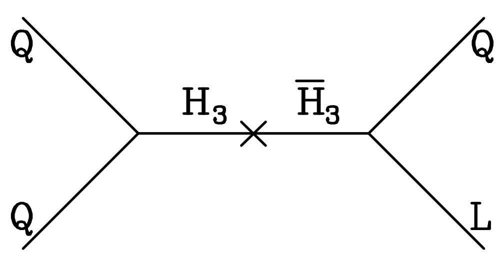

Figure 2: Dimension-five proton decay operator obtained from the FFh, $F\bar f\bar h$, and $H_3\bar H_3$ interactions which, when suitably dressed by squark and gaugino loops, mediates the decay channel $p \to \bar\nu K^+$. This channel is negligible in flipped SU(5) because of the absence of the $H_3\bar H_3$ mixing term.

1.3 Prediction for $\alpha_s(M_Z)$

In flipped SU(5), there is a first unification scale M_{32} at which the SU(3) and SU(2) gauge couplings become equal, given to lowest order by [11]

$$\frac{1}{\alpha_2} - \frac{1}{\alpha_5} = \frac{b_2}{2\pi} \ln \frac{M_{32}}{M_Z} , \tag{12}$$

$$\frac{1}{\alpha_3} - \frac{1}{\alpha_5} = \frac{b_3}{2\pi} \ln \frac{M_{32}}{M_Z} , \tag{13}$$

where $\alpha_2 = \alpha / \sin^2 \theta_W$, $\alpha_3 = \alpha_s(M_Z)$, and the one-loop beta functions are $b_2 = +1$, $b_3 = -3$. On the other hand, the hypercharge gauge coupling $\alpha_Y = \frac{5}{3}(\alpha / \cos^2 \theta_W)$ evolves in general to a different value α_1' at the scale M_{32}:

$$\frac{1}{\alpha_Y} - \frac{1}{\alpha_1'} = \frac{b_Y}{2\pi} \ln \frac{M_{32}}{M_Z} , \tag{14}$$

with $b_Y = \frac{33}{5}$. Above the scale M_{32} the gauge group is SU(5)×U(1), with the U(1) gauge coupling α_1 related to α_1' and the SU(5) gauge coupling (α_5) by

$$\frac{25}{\alpha_1'} = \frac{1}{\alpha_5} + \frac{24}{\alpha_1} . \tag{15}$$

The SU(5) and U(1) gauge couplings continue to evolve above the scale M_{32}, eventually becoming equal at a higher scale M_{51}. The consistency condition that $M_{51} \geq M_{32}$ implies $\alpha_1' \leq \alpha_5(M_{32})$. The maximum possible value of M_{32} is obtained when $\alpha_1' = \alpha_5(M_{32})$ and is given by

$$\frac{1}{\alpha_Y} - \frac{1}{\alpha_5} = \frac{b_Y}{2\pi} \ln \frac{M_{32}^{\max}}{M_Z} . \tag{16}$$

Solving the above equations for the value of $\alpha_s(M_Z)$ we obtain [12]

$$\alpha_s(M_Z) = \frac{\frac{7}{3}\,\alpha}{5 \sin^2 \theta_W - 1 + \frac{11}{2\pi}\,\alpha \ln(M_{32}^{\max}/M_{32})} . \tag{17}$$

Note that since in minimal SU(5) $M_{32} = M_{32}^{\max}$, we automatically obtain

$$\alpha_s(M_Z)^{\text{Flipped SU(5)}} < \alpha_s(M_Z)^{\text{SU(5)}} . \tag{18}$$

The next-to-leading order corrections to Eq. (17) are obtained by the substitution

$$\sin^2 \theta_W \rightarrow \sin^2 \theta_W - \delta_{\text{2loop}} - \delta_{\text{light}} - \delta_{\text{heavy}} , \tag{19}$$

where δ_{2loop} accounts for the two-loop contributions to the RGEs, δ_{light} accounts for the light SUSY thresholds, and δ_{heavy} accounts for the GUT thresholds. In both minimal SU(5) and flipped SU(5) one finds $\delta_{\text{2loop}} \approx 0.0030$ and $\delta_{\text{light}} \gtrsim 0$. These effects go the

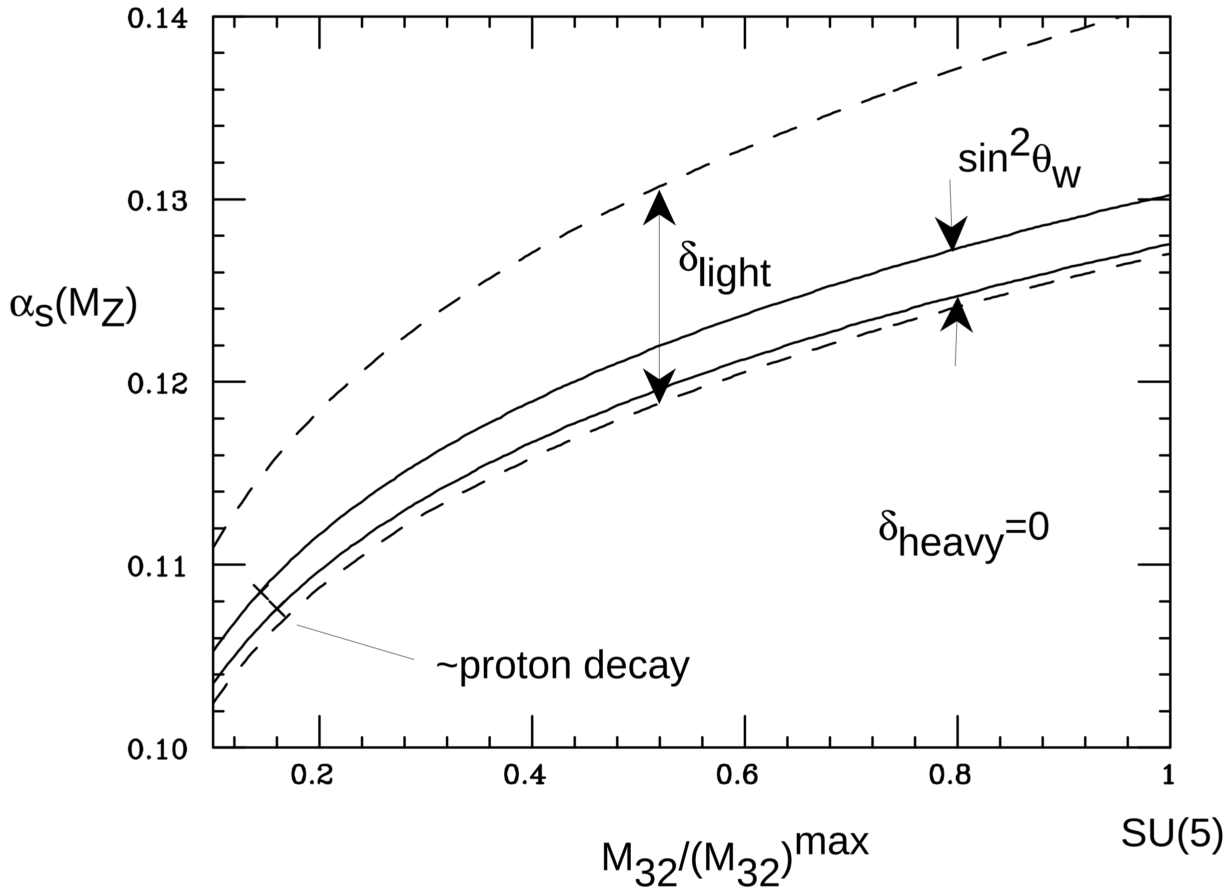

Figure 3: Prediction for $\alpha_s(M_Z)$ in flipped SU(5) as a function of $M_{32}/M_{32}^{\max}$. The solid lines represent the range of predictions for $\sin^2\theta_W = 0.23143 \pm 0.00028$ with no threshold corrections ($\delta_{\rm light} = \delta_{\rm heavy} = 0$). The dashed lines represent the excursion obtained by scanning over sparticle masses ($\delta_{\rm light} \neq 0$) below 1 TeV.

wrong way for minimal SU(5), and lead to the lower bound $\alpha_s^{\rm SU(5)}(M_Z) > 0.123$ [4]. There is no help from $\delta_{\rm heavy}$ as it has to be positive to minimize the proton lifetime via dimension-five operators. In the case of flipped SU(5), the reduced unification scale allows a lower prediction for $\alpha_s(M_Z)$, and one could even get further help from $\delta_{\rm heavy}$ [12]

$$\delta_{\rm heavy} = \frac{\alpha}{20\pi}\left[-6\ln\frac{M_{32}}{M_{H_3}} - 6\ln\frac{M_{32}}{M_{\bar{H}_3}} + 4\ln\frac{M_{32}}{M_V}\right] . \tag{20}$$

Since there is no problem with proton decay, the $H_3, \bar{H}_3$ masses can be lighter than M_{32} and $\delta_{\rm heavy} > 0$ is perfectly acceptable. In Fig. 3 we show the prediction for $\alpha_s(M_Z)$ in flipped SU(5) as a function of the ratio $M_{32}/M_{32}^{\max}$. We note that the minimal SU(5) prediction is obtained when $M_{32}/M_{32}^{\max} = 1$. Scanning over the possible range of sparticle masses described above, we find significant variations in the predictions for $\alpha_s(M_Z)$, indicated by the dashed lines in Fig. 3, usually towards higher values

(*i.e.*, $\delta_{\rm light} > 0$). Equations (19) and (17) show that including the effects of $\delta_{\rm heavy}$ simply amounts to a re-scaling of the $M_{32}/M_{32}^{\rm max}$ axis on Fig. 3, *i.e.*,

$$\frac{M_{32}}{M_{32}^{\rm max}} \to \frac{M_{32}}{M_{32}^{\rm max}}\ e^{-10\pi\,\delta_{\rm heavy}/11\alpha}\ . \tag{21}$$

Decreasing the unification scale enhances the proton decay rate. One obtains [12]

$$\tau(p\to e^+\pi^0) \approx 1.5\times 10^{33}\left(\frac{M_{32}}{10^{15}\,{\rm GeV}}\right)^4\left(\frac{0.042}{\alpha_5}\right)^2, \tag{22}$$

In Fig. 4 we plot $\alpha_s(M_Z)$ versus $\tau(p\to e^+\pi^0)$. The present experimental lower bound $\tau(p\to e^+\pi^0)^{\rm exp} > 5.5\times 10^{32}\,$y allows values of $\alpha_s(M_Z)$ as low as 0.108. Moreover, for values of $\alpha_s(M_Z) \lesssim 0.114$, the mode $p\to e^+\pi^0$ should be observable at SuperKamiokande in flipped SU(5), whereas in minimal supersymmetric SU(5) the dominant mode is expected to be $p\to\bar\nu K^+$.

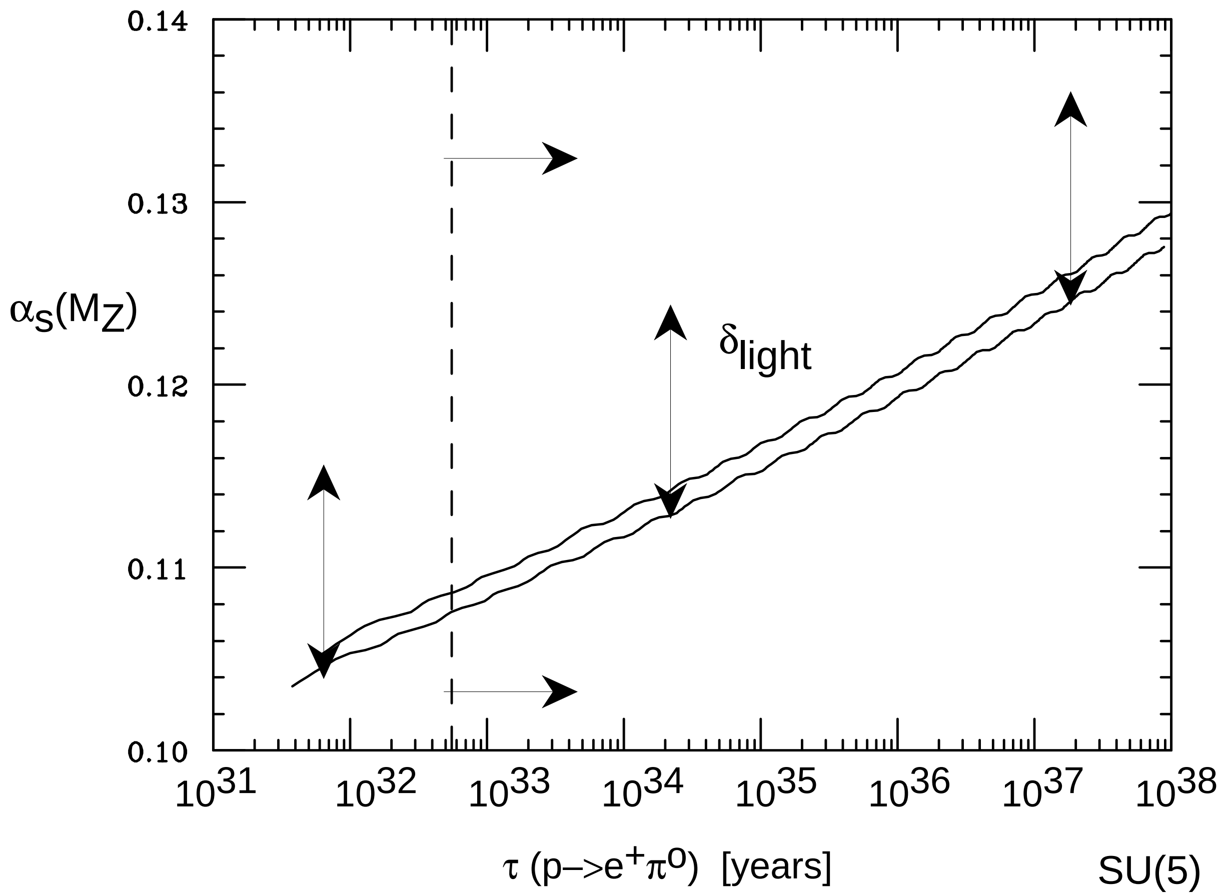

Figure 4: Prediction for $\alpha_s(M_Z)$ in flipped SU(5) as a function of the proton lifetime into the mode $e^+\pi^0$. The present experimental lower bound is indicated by the dashed line. Vertical arrows indicate the effect of $\delta_{\rm light}$.

2 Stringy Flipped SU(5)

The flipped SU(5) GUT model described in the previous section has many nice features that make it a very appealing model of physics at very high energies. However, it leaves several unanswered questions

- What are the particle masses (e.g., m_t)?
- What are superparticle masses (e.g., $m_{\tilde{q}}$)?
- What about unification with gravity?

Making supersymmetry a local symmetry (*i.e.*, supergravity) has several advantages. For instance, the superparticle masses can be parametrized in terms of much fewer parameters. A popular scenario assumes universality of these supersymmetry-breaking mass terms, resulting in a model with only four parameters. Still, gravity is only incorporated as an effective theory below the Planck mass. But even supergravity is not enough, as it leaves its own set of unanswered questions

- What are the new supergravity functions?
- What about gravity <u>at</u> the Planck scale?
- What about the particle masses?

These appear rather daunting questions, especially the last one

> "No theory can predict any particle mass" – Not!

In fact, these three questions can be answered in the context of string models although, because of the large number of possible models, the answers at the moment are not unique. String theorists consider this to be no prediction at all. However in practice, even though there exist many possible models, and even some semi-realistic ones (like stringy flipped SU(5)), one can hardly claim that the one and only string model has been found. Thus, the search for the best string model appears to be a perfectly legitimate and physically reasonable pursuit. (It is also possible that somehow there is not a unique vacuum (or model) of string.)

2.1 String derivation of Flipped SU(5)

Flipped SU(5) was originally obtained as a string model in its so-called "revamped" version in 1989 [13]. The phenomenological properties of this model were then extensively explored from first-principles string calculations [14, 15]. With time it became apparent that it was difficult to achieve unification of the gauge couplings near the string scale, as such scenario required additional specific intermediate-scale matter representations with Standard Model quantum numbers [16], which were not available. In 1992 a systematic exploration of a large class of string flipped SU(5) models

[17] led to a new model with a one extra pair of ($\mathbf{10}$,$\overline{\mathbf{10}}$) representations that in principle allowed string unification. Here we outline the derivation and main properties of this latest and more realistic flipped SU(5) string model [17].

As is well known, in free-fermionic string model building one must provide a set of basis vectors of boundary conditions for the two-dimensional world-sheet fermions as they traverse the one-loop (torus) world-sheet. Our choice is:

$$
\begin{array}{rll}
\mathbf{1} & = (1\ 111\ 111\ 111\ 111\ 111\ 111 & :\ 111111\ 111111\ 11111\ 111\ 1_8) \\
S & = (1\ 100\ 100\ 100\ 100\ 100\ 100 & :\ 000000\ 000000\ 00000\ 000\ 0_8) \\
b_1 & = (1\ 100\ 100\ 010\ 010\ 010\ 010 & :\ 001111\ 000000\ 11111\ 100\ 0_8) \\
b_2 & = (1\ 010\ 010\ 100\ 100\ 001\ 001 & :\ 110000\ 000011\ 11111\ 010\ 0_8) \\
b_3 & = (1\ 001\ 001\ 001\ 001\ 100\ 100 & :\ 000000\ 111100\ 11111\ 001\ 0_8) \\
b_4 & = (1\ 100\ 100\ 010\ 001\ 001\ 010 & :\ 001001\ 000110\ 11111\ 100\ 0_8) \\
b_5 & = (1\ 001\ 010\ 100\ 100\ 001\ 010 & :\ 010001\ 100010\ 11111\ 010\ 0_8) \\
\alpha & = (0\ 000\ 000\ 000\ 000\ 000\ 011 & :\ 000001\ 011001\ \frac{1}{2}\frac{1}{2}\frac{1}{2}\frac{1}{2}\frac{1}{2}\ \frac{1}{2}\frac{1}{2}\frac{1}{2}\ A)
\end{array}
$$

One also needs to specify a matrix of GSO projections which determine which states in the spectrum are allowed by the modular invariance properties of the two-dimensional world-sheet

$$
k = \begin{pmatrix}
2 & 1 & 2 & 2 & 2 & 2 & 2 & 1 \\
1 & 1 & 1 & 1 & 1 & 1 & 1 & 4 \\
2 & 2 & 2 & 2 & 2 & 2 & 2 & 1 \\
2 & 2 & 2 & 2 & 2 & 2 & 2 & 1 \\
2 & 2 & 2 & 2 & 2 & 2 & 2 & 4 \\
2 & 2 & 2 & 2 & 2 & 2 & 2 & 1 \\
2 & 2 & 2 & 2 & 2 & 2 & 2 & 3 \\
2 & 1 & 1 & 1 & 1 & 1 & 2 & 3
\end{pmatrix} \tag{23}
$$

The well-known rules of free-fermionic model-building can then be employed to deduce a series of properties of the resulting model, such as the gauge group, the particle spectrum, and all the interactions in the model. The gauge group $G = G_{\rm observable} \times G_{\rm hidden} \times U(1)^5$ contains an observable gauge group $G_{\rm observable}$=SU(5)×U(1) and a hidden gauge group $G_{\rm hidden}$=SO(10)×SU(4). The particle spectrum divides itself into a finite number of "massless" particles and an infinite tower of Planck mass particles. The massless fields include:

- Observable Sector:

$$
\begin{array}{lll}
F^{\{0,1,2,3,4\}}[10] & \bar{f}^{\{2,3,5\}}[\bar{5}] & \ell^{c\{2,3,5\}}[1] \\
\bar{F}^{\{4,5\}}[\overline{10}] & & \\
h^{\{1,2,3,45\}}[5] & \bar{h}^{\{1,2,3,45\}}[\bar{5}] &
\end{array} \tag{24}
$$

- Singlet fields: 20 charged, 4 neutral under all gauge symmetries ($\Phi_{0,1,3,5}$).

- Hidden Sector:

$$\begin{array}{ll} \mathrm{T}^{\{1,2,3\}} & [10] \text{ of } \mathrm{SO}(10) \\ \mathrm{D}^{\{1,2,3,4,5,6,7\}} & [6] \text{ of } \mathrm{SU}(4) \\ \widetilde{\mathrm{F}}^{\{1,2,3,4,5,6\}} & [4] \text{ of } \mathrm{SU}(4) \\ \widetilde{\widetilde{\mathrm{F}}}^{\{1,2,3,4,5,6\}} & [\bar{4}] \text{ of } \mathrm{SU}(4) \end{array} \tag{25}$$

The $\widetilde{F}_i, \widetilde{\widetilde{F}}_j$ fields carry $\pm 1/2$ electric charges and exist only confined in hadron-like *cryptons* [18]. This is crucial phenomenological property which is achieved naturally in stringy flipped SU(5) [19].[2]

The cubic superpotential can be easily determined to be

$$\begin{array}{rcl} W_3 = g\sqrt{2} & \Big\{ & F_0 F_1 h_1 + F_2 F_2 h_2 + F_4 F_4 h_1 + F_4 \bar{f}_5 \bar{h}_{45} + F_3 \bar{f}_3 \bar{h}_3 \\ & + & \bar{f}_2 l_2^c h_2 + \bar{f}_5 l_5^c h_2 \\ & + & \frac{1}{\sqrt{2}} F_4 \bar{F}_5 \phi_3 + \frac{1}{2} F_4 \bar{F}_4 \Phi_0 + \bar{F}_4 \bar{F}_4 \bar{h}_1 + \bar{F}_5 \bar{F}_5 \bar{h}_2 \\ & + & (h_1 \bar{h}_2 \Phi_{12} + h_2 \bar{h}_3 \Phi_{23} + h_3 \bar{h}_1 \Phi_{31} + h_3 \bar{h}_{45} \bar{\phi}_{45} + \text{h.c.}) \\ & + & \frac{1}{2}(\phi_{45} \bar{\phi}_{45} + \phi^+ \bar{\phi}^+ + \phi^- \bar{\phi}^- + \phi_i \bar{\phi}_i + h_{45} \bar{h}_{45}) \Phi_3 \\ & + & (\eta_1 \bar{\eta}_2 + \bar{\eta}_1 \eta_2) \Phi_0 + (\phi_3 \bar{\phi}_4 + \bar{\phi}_3 \phi_4) \Phi_5 \\ & + & (\Phi_{12} \Phi_{23} \Phi_{31} + \Phi_{12} \phi^+ \phi^- + \Phi_{12} \phi_i \phi_i + \text{h.c.}) \\ & + & T_1 T_1 \Phi_{31} + T_3 T_3 \Phi_{31} \\ & + & D_6 D_6 \Phi_{23} + D_1 D_2 \bar{\Phi}_{23} + D_5 D_5 \bar{\Phi}_{23} + D_7 D_7 \bar{\Phi}_{31} \\ & + & D_3 D_3 \Phi_{31} + \frac{1}{2} D_5 D_6 \Phi_0 + \frac{1}{\sqrt{2}} D_5 D_7 \bar{\phi}_3 \\ & + & \widetilde{F}_4 \widetilde{\widetilde{F}}_6 \bar{\Phi}_{12} + \frac{1}{2} F_3 \widetilde{\widetilde{F}}_4 \Phi_0 + \frac{1}{2} F_2 \widetilde{\widetilde{F}}_5 \Phi_3 + \widetilde{F}_6 \widetilde{\widetilde{F}}_4 \phi^+ \\ & + & \frac{1}{\sqrt{2}} \widetilde{F}_5 \widetilde{\widetilde{F}}_4 \phi_4 + \widetilde{F}_1 \widetilde{\widetilde{F}}_2 D_5 + \widetilde{F}_2 \widetilde{\widetilde{F}}_4 l_2^c \Big\} \end{array}$$

Higher-order terms in the superpotential are suppressed by powers of the string scale, but are also calculable. The quartic superpotential terms are

$$\begin{array}{rcl} W_4 & = & F_2 \bar{f}_2 \bar{h}_{45} \bar{\phi}_4 + F_3 \bar{F}_4 D_4 D_6 + F_3 \bar{F}_5 D_4 D_7 \\ & + & l_3^c \widetilde{\widetilde{F}}_3 \widetilde{\widetilde{F}}_6 D_7 + l_5^c \widetilde{F}_2 \widetilde{\widetilde{F}}_3 \bar{\phi}_3 + \widetilde{F}_1 \widetilde{\widetilde{F}}_3 (\phi^+ \bar{\phi}_3 + \bar{\phi}^- \phi_3) \\ & + & \widetilde{\widetilde{F}}_3 \widetilde{\widetilde{F}}_5 D_7 \bar{\phi}^- + \widetilde{F}_2 \widetilde{F}_5 D_3 \phi^- + \widetilde{F}_2 \widetilde{F}_6 D_3 \phi_4 + \widetilde{F}_5 \widetilde{\widetilde{F}}_1 D_2 D_7 \\ & + & \widetilde{F}_5 \widetilde{\widetilde{F}}_2 D_1 D_7 + \widetilde{F}_3 \widetilde{\widetilde{F}}_3 D_3 D_6 + \widetilde{F}_4 \widetilde{\widetilde{F}}_3 D_4 D_7 + \widetilde{F}_5 \widetilde{\widetilde{F}}_4 D_5 D_7. \end{array}$$

where the coefficients are generically $\mathcal{O}(1)$ and can be calculated explicitly [14]. Traditionally the above results have been used to study the phenomenology of the models, aided by a set of non-zero vacuum expectation values for many of the singlets fields, as required by the flatness conditions in the presence of an anomalous $\mathrm{U_A}(1)$ (see below).

[2] This is to be contrasted with other fermionic models where the confinement of such particles may be quite difficult to achieve [20], and yet further models where fractionally charged particles appear in the observable sector [21] and thus confinement is not even a possibility.

2.2 Recent developments

Not until recently has the Kähler potential of realistic fermionic models been studied in detail [22]. This second phase of string model building is essential to the study of the supersymmetry-breaking spectrum that arises, and depends crucially on the proper identification of the moduli fields in free-fermionic models.

2.2.1 The Kähler potential

As a first step, one notes that the fields in the model organize themselves into three sets of untwisted and twisted fields. In this model one finds

$$
\begin{array}{lcc}
\underline{\text{Set}} & \underline{\text{Untwisted}}\ [\alpha_i^{(I)} \equiv U^{(I)}] & \underline{\text{Twisted}}\ [\beta_i^{(I)} \equiv T^{(I)}] \\
\text{First} & \Phi_0, \Phi_1, \Phi_{23}, \bar{\Phi}_{23} & F_0, F_1, F_4, \bar{F}_4 \\
 & h_1, \bar{h}_1 & \widetilde{F}_3, \widetilde{\bar{F}}_{1,2,4}, D_{1,2,5,6} \\
\text{Second} & \eta_1, \bar{\eta}_1, \Phi_{31}, \bar{\Phi}_{31} & F_2, \bar{f}_2, l_2^c; \bar{F}_5, \bar{f}_5, l_5^c \\
 & h_2, \bar{h}_2 & \widetilde{F}_{1,5,6}, \widetilde{\bar{F}}_3, D_{3,7}, T_{1,3} \\
\text{Third} & \Phi_3, \Phi_5, \Phi_{12}, \bar{\Phi}_{12} & F_3, \bar{f}_3, l_3^c; h_{45}, \bar{h}_{45} \\
 & \eta_2, \bar{\eta}_2 & \phi_{45}, \bar{\phi}_{45}, \phi^+, \bar{\phi}^+, \phi^-, \bar{\phi}^-, \phi_{3,4}, \bar{\phi}_{3,4} \\
 & h_3, \bar{h}_3 & \widetilde{F}_{2,4}, \widetilde{\bar{F}}_{5,6}, D_4, T_2
\end{array}
\tag{26}
$$

These sets are easily deduced by studying some left-moving (*i.e.*, world-sheet supersymmetric) charge assignments. The Kähler potential can then be calculated. The crucial observation is that the model possess only one modulus field (besides the dilaton S) which corresponds to the all-neutral untwisted singlet field Φ_1, which belongs to the first set. One finds [23]

$$
\begin{aligned}
K \;=\; & -\ln(S+\bar{S}) - \ln\left[(\tau+\bar{\tau})^2 - \sum_i^{n_{U_1}} [\alpha_i^{(1)} + \bar{\alpha}_i^{(1)}]^2\right] \\
& + \sum_i^{n_{U_2}} \alpha_i^{(2)} \bar{\alpha}_i^{(2)} + \sum_i^{n_{U_3}} \alpha_i^{(3)} \bar{\alpha}_i^{(3)} + \sum_i^{n_{T_1}} \beta_i^{(1)} \bar{\beta}_i^{(1)} \\
& + \frac{1}{\left[(\tau+\bar{\tau})^2 - \sum_i^{n_{U_1}} [\alpha_i^{(1)} + \bar{\alpha}_i^{(1)}]^2\right]^{1/2}} \left(\sum_i^{n_{T_2}} \beta_i^{(2)} \bar{\beta}_i^{(2)} + \sum_i^{n_{T_3}} \beta_i^{(3)} \bar{\beta}_i^{(3)} \right)
\end{aligned}
\tag{27}
$$

Note the special role played by the modulus field (τ) and the other fields in the first unstwisted set. The gravitino mass is given by

$$
m_{3/2}^2 = e^{\langle K \rangle} \langle |W| \rangle^2 = \frac{\langle |W| \rangle^2}{\langle (S+\bar{S})(\tau+\bar{\tau})^2 \rangle}
\tag{28}
$$

where the sole source of supersymmetry breaking is $\langle W \rangle \neq 0$. In what follows we assume that $\langle W \rangle$ is a constant independent of the moduli fields. Moreover, as it stands the gravitino mass is undetermined, since it depends on the flat directions S

and τ. In what follows we assume that these are somehow fixed to $\langle S\rangle \sim \langle\tau\rangle \sim 1$ (in Planck units), as motivated from S and T duality considerations. To achieve a sufficiently small value of $m_{3/2}$ then $\langle W\rangle$ would have to be suppressed, as occurs for instance in gaugino condensation models. From the low-energy effective theory point of view, $\langle S\rangle$ and $\langle\tau\rangle$ could be determined dynamically via the no-scale mechanism [24]. This mechanism is destabilized by one-loop quadratically divergent contributions to the scalar potential, but fortunately in our model (as discussed in Sec. 2.2.4) these can be made to vanish in a suitably chosen vacuum.

2.2.2 Properties of the Kähler potential [23]

The most basic property of the Kähler potential in Eq. (27) is that the tree-level vacuum energy vanishes, *i.e.*,

$$V_0 = m_{3/2}^2 \left(1[S] + 2[\tau] - 3\right) = 0 \tag{29}$$

where the contributions from the two moduli are indicated. Next we need to worry about the quadratically-divergent one-loop correction [25, 26]

$$\frac{1}{32\pi^2} \mathrm{Str}\, \mathcal{M}^2 M_{\mathrm{Pl}}^2 \equiv \frac{1}{16\pi^2}\, Q_0\, m_{3/2}^2\, M_{\mathrm{Pl}}^2 \tag{30}$$

The quantity Q_0 can be straightforwardly calculated from the Kähler potential

$$Q_0 = n_{U_2} + n_{U_3} + n_{T_1} - n_{U_1} - d_f - 3 = 4 \tag{31}$$

where the $n_{U_1} = 13, n_{U_2} = 14, n_{U_3} = 16$ and $n_{T_1} = 80, n_{T_2} = 80, n_{T_3} = 68$ represent the numbers of fields in each of the untwisted and twisted sets in Eq. (26), and $d_f = 90$ is the dimension of the gauge group. We note that this lowest-order calculation of Q_0 yields a "small" enough value, which can be further shifted towards zero in the presence of an anomalous $\mathrm{U_A}(1)$ factor.

Turning to the supersymmetry-breaking parameters, the goldstino field is

$$\tilde{\eta} = \frac{S + \sqrt{2}\,\tau}{\sqrt{3}} \tag{32}$$

whereas the scalar masses are given by the following multiples of $m_{3/2}$

$$\begin{array}{llllll} U^{(1)}: & \alpha^{(1)} & 0 & \qquad T^{(1)}: & \beta_i^{(1)} & 1 \\ U^{(2)}: & \alpha_i^{(2)} & 1 & \qquad T^{(2)}: & \beta_i^{(2)} & 0 \\ U^{(3)}: & \alpha_i^{(3)} & 1 & \qquad T^{(3)}: & \beta_i^{(3)} & 0 \end{array} \tag{33}$$

The trilinear scalar couplings are universal

$$A = m_{3/2} \tag{34}$$

and so are the (tree-level) gaugino masses

$$m_{1/2} = m_{3/2} \tag{35}$$

2.2.3 Flatness constraints

Keeping supersymmetry unbroken at the string scale requires the usual F- and D-flatness conditions to be satisfied. The solutions to these conditions become non-trivial in the presence of an anomalous $\mathrm{U_A}(1)$ factor in the gauge group, as is the case of all free-fermionic models known to date. F-flatness is given by the usual condition

$$\left\langle \frac{\partial W}{\partial \phi_i} \right\rangle = 0 \tag{36}$$

where the ϕ_i run over all fields in the model. D-flatness involves the usual conditions for the non-anomalous U(1) factors in the gauge group

$$\langle D_a \rangle = \sum_i q_a^i \langle \phi \rangle^2 = 0 \tag{37}$$

and the modified relation for $\mathrm{U_A}(1)$

$$\langle D_A \rangle = \sum_i q_A^i \langle \phi \rangle^2 + \epsilon = 0\,, \quad \epsilon = \frac{g^2 M^2}{192\pi^2}\,\mathrm{Tr}\,\mathrm{U_A} > 0 \tag{38}$$

In this expression the fundamental scale is $M \approx 10^{18}\,\mathrm{GeV}$, from which we can obtain (in this model) the new scale $\sqrt{\epsilon}$ [17]

$$\sqrt{\epsilon} = 3.7 \times g \times 10^{17}\,\mathrm{GeV} \sim \frac{M}{10} \tag{39}$$

The scalar fields charged under $\mathrm{U_A}(1)$ get vevs $\mathcal{O}(\sqrt{\epsilon})$, such that all the F- and D-flatness constraints are satisfied. This procedure amounts to a shift in the vacuum, from the original one built within fermionic strings (where all vevs vanish), to a "nearby" one where supersymmetry is unbroken (and the vevs are $\mathcal{O}(\sqrt{\epsilon})$). The "shifted" vacuum is still a consistent string solution, but which is not directly reachable from within fermionic strings. Nonetheless, in models with $\mathrm{U_A}(1)$, starting from the original vacuum, one can sample a large number of nearby vacua by the deformations that preserve flatness, which parametrize a multi-dimensional space of allowed vevs.

The usual F- and D-flatness conditions above are sufficient as long as one does not consider supersymmetry-breaking effects that may give masses to some of the scalars. This is certainly the case in the model at hand, as seen in Eq. (33). If one were to consider giving a vev to a scalar field that acquires a supersymmetry breaking mass, in the shifted vacuum the vacuum energy will not vanish anymore. Thus, one restricts the choice of fields to be shifted to those that do not acquire supersymmetry-breaking masses [23]. This is a new constraint in string model-building, which leads to more restrictive sets of allowed solutions for the shifted vevs.

2.2.4 Shifting Q [26]

In the process of going to the shifted vacuum, the supersymmetry-breaking masses that contribute to $\mathrm{Str}\,\mathcal{M}^2$ may be shifted as well, leading to a shift in Q_0. One can show that in our model the ensuing shift always lowers the value of Q_0 (more details are given below)

$$Q = Q_0 - (\text{specific sum of vevs squared}) \tag{40}$$

One is then led to the question: Can one find a set of vevs such that F-flatness is Ok, D-flatness in the presence of $\mathrm{U_A}(1)$ is Ok, and Q is shifted to zero? The answer is yes, several such solutions have been found explicitly [28]. Moreover, the explicit solutions found have been shown to be rather typical, demonstrating the wide applicability of this mechanism [27].

We should remark that this mechanism to shift Q towards zero works only if $V_0 = 0$ and Q_0 is "small" to begin with. In our specific example Q_0 must be positive too. If $V_0 \neq 0$ then in the expression for Q_0 (c.f. Eq. (31)) the relative signs of the various terms do not cooperate and invariably one obtains large values of Q_0. In fact, the model discussed here is the only known free-fermionic model where $V_0 = 0$, and it is the only known free-fermionic model where Q_0 is "small" and therefore can be shifted to zero by our mechanism. In all other known models (see Table 3) $V_0 \neq 0$ and $|Q_0| = \mathcal{O}(100)$. It is also worth remarking that being able to *calculate* V_0 and Q is in itself a non-trivial matter, which has not really been done in many models so far, perhaps because of the yet-to-be-calculated form of the Kähler potential.

Table 3: Compilation of V_0 and Q_0 values in known fermionic string models.

Model	V_0	Q_0
PZ [29]	0	-272
ALR [30]	3	-269
F [20]	-2	168
AEHN [13]	1	-83
LNY [17]	0	4

Having demonstrated the viability of models with vanishing values of Q, the question remains whether higher-order contributions to the scalar potential may generate new quadratically divergent contributions, as pointed out in Ref. [31]. In other words, is there any reason to believe that, taking into account the full string theory, these higher-order quadratic divergences somehow vanish automatically, as long as Q vanishes? A would-be analogous situation occurs with the string loop corrections to the gauge kinetic function, which vanish at two and higher loops. In this case a modular anomaly is at play. In our case we would like to conjecture that the anomalous $\mathrm{U_A}(1)$ plays the role of the anomaly at hand and, as with all such anomalies, as long as one takes care of their appearance at one-loop order, higher orders are dealt with automatically. Detailed string calculations appear needed to verify or disprove such conjecture.

2.2.5 Specific example of $\mathrm{U_A}(1)$ and Q cancellation [27]

As discussed above, in order to demonstrate the possibility of achieving F- and D-flatness, in the presence of supersymmetry-breaking masses, we need to impose

$$\begin{aligned} \alpha^{(2)} &: \left\langle \eta_1, \bar{\eta}_1, \Phi_{31}, \bar{\Phi}_{31} \right\rangle = 0 \\ \alpha^{(3)} &: \left\langle \eta_2, \bar{\eta}_2, \Phi_{12}, \bar{\Phi}_{12}, \Phi_3, \Phi_5 \right\rangle = 0 \\ \beta^{(1)} &: \langle \nu_0^c, \nu_1^c, \nu_4^c, \nu_{\bar{4}}^c \rangle = 0 \end{aligned} \tag{41}$$

so that the vacuum energy remains zero in the shifted vacuum. The usual F- and D-flatness conditions become more restrictive and amount to the following set of constraints (numbers in square brackets refer to the original set of constraints in Ref. [17], where the new constraints in Eq. (41) were not imposed)

$$\begin{aligned}
&[6.7] \quad \begin{cases} x_{45} = \frac{1}{15}\epsilon - \frac{1}{2}V_3^2 \\ x_+ - x_- = \frac{1}{15}\epsilon + \frac{1}{2}V_3^2 \\ x_{23} = -\frac{1}{5}\epsilon \\ x_3 + x_4 + 2x_+ = \frac{2}{5}\epsilon + V_3^2 \end{cases} \\
&[6.8] \quad \left\{ V^2 = V_2^2 + V_3^2 = \bar{V}_5^2 = \bar{V}^2 \right. \\
&[6.9] \quad \begin{cases} \langle \phi^+\phi^- + \phi_4\phi_4 \rangle = 0 \\ \langle \bar{\phi}^+\bar{\phi}^- + \bar{\phi}_3\bar{\phi}_3 + \bar{\phi}_4\bar{\phi}_4 \rangle = 0 \\ \langle \phi_{45}\bar{\phi}_{45} + \phi_4\bar{\phi}_4 + \phi^+\bar{\phi}^+ + \phi^-\bar{\phi}^- \rangle = 0 \\ \langle \bar{\phi}_3\phi_4 \rangle = 0 \\ \langle \phi_3 \rangle = 0 \end{cases}
\end{aligned}$$

where $x_i = |\langle\varphi_i\rangle|^2 - |\langle\bar{\varphi}_i\rangle|^2$ and $V_i = \langle\nu_i^c\rangle$. It also follows that Q is shifted in the shifted vacuum [27]

$$Q = 4 - \tfrac{1}{3}(14X + 5Y) \tag{42}$$

where

$$X = \sum_{\alpha^{(1)}} \frac{\langle \alpha^{(1)} + \bar{\alpha}^{(1)} \rangle^2}{\langle \tau + \bar{\tau} \rangle^2}\ , \qquad Y = \sum_{\beta^{(2)}} \frac{\langle \beta^{(2)}\bar{\beta}^{(2)} \rangle}{\langle \tau + \bar{\tau} \rangle} + \sum_{\beta^{(3)}} \frac{\langle \beta^{(3)}\bar{\beta}^{(3)} \rangle}{\langle \tau + \bar{\tau} \rangle} \tag{43}$$

Note that $X > 0$ and $Y > 0$, and therefore Q is always shifted towards zero. It is important to note that this expression for Q in terms of the X, Y variables has been obtained in the reasonable approximation that $\langle\tau\rangle \sim 1$ (as assumed after Eq. (28)) and $\langle\phi\rangle \ll \langle\tau\rangle$, as would be expected from anomalous $\mathrm{U_A}(1)$ considerations. As such, large shifts in Q are by construction not possible. Also, the duality properties of the Kähler potential may not manifest in the expression for Q because of the approximation made. The possible non-zero contributors to X, Y are

$$\begin{aligned} \alpha^{(1)} &: \left\langle \Phi_0, \Phi_{23}, \bar{\Phi}_{23} \right\rangle \\ \beta^{(2)} &: V_2, \bar{V}_5 \\ \beta^{(3)} &: V_3, \left\langle \phi_{45}, \bar{\phi}_{45}, \phi^+, \bar{\phi}^+, \phi^-, \bar{\phi}^-, \phi_{3,4}, \bar{\phi}_{3,4} \right\rangle \end{aligned} \tag{44}$$

Table 4: Sample solutions to the F- and D-flatness constraints, respecting the new constraints from $V_0 = 0$ in the shifted vacuum. All vevs in units of $\epsilon^{1/2}$.

$\widehat{\phi}_{45}$	$\widehat{\overline{\phi}}_{45}$	$\widehat{\phi}^{+}$	$\widehat{\overline{\phi}}^{+}$	$\widehat{\phi}^{-}$	$\widehat{\overline{\phi}}^{-}$	$\widehat{\phi}_4$	$\widehat{\overline{\phi}}_4$	$\widehat{V}_3$	$\sum \widehat{\beta}^2$
0.836	-0.796	0.531	0.368	-0.539	-0.472	0.535	0.417	0.060	2.73
0.999	-1.072	0.942	0.741	-0.386	-0.279	0.603	0.455	0.659	5.25
0.275	-0.132	0.401	0.089	-0.256	-0.008	0.320	0.026	0.130	0.46
0.589	-0.702	0.675	0.341	-0.298	-0.213	0.448	0.269	0.650	2.66
0.735	-0.910	0.892	0.559	-0.265	-0.150	0.486	0.290	0.842	4.31
0.835	-0.837	0.376	0.051	-0.750	-0.762	0.531	0.197	0.374	3.28
0.579	-0.545	0.401	0.074	-0.467	-0.416	0.433	0.175	0.236	1.52
0.927	-0.987	0.829	0.613	-0.435	-0.344	0.600	0.459	0.603	4.50

Scanning over the multi-dimensional parameter space of vevs one can find solutions to the F- and D-flatness conditions (respecting the new constraints in Eq. (41)). Sample values are shown in Table 4. One can then see whether these solutions also lead to $Q = 0$. This is the case for several of them (the ones with the larger values of $\sum \widehat{\beta}^2$), as Q also depends on vevs (in the X contribution) which are not too constrained by the flatness conditions, and thus can be adjusted to yield $Q = 0$.

2.3 Particle masses

Superparticle masses receive contributions from two sources: from the masses of their Standard Model partners and from supersymmetry breaking. Experimentally we now know that the latter contribution dominates (except perhaps in the case of a light top-squark). These supersymmetry-breaking contributions have been given above (see Eq. (33)) and scale with $m_{3/2}$. On the other hand, particle masses (such as quarks and leptons) can be determined from three inputs:

- A mass scale, which must be M_Z as masses are protected (*i.e.*, vanish) above the scale of SU(2)×U(1) breaking
- A Yukawa coupling, which is a pure number calculable only in string models, where one typically obtains

  $$\lambda \sim g \quad \text{(gauge coupling)} \tag{45}$$

  Note that GUTs can <u>relate</u> but not predict Yukawa couplings.
- A dynamical coefficient coming from the evolution of the Yukawa couplings from the string scale down to the electroweak scale ($\lambda(M_{Pl}) \to \lambda(M_Z)$), and from low-energy mixing effects (*e.g.*, $\tan\beta$).

Given this scenario one would then expect $m_q \sim \lambda M_Z$, or in somebody else's words

"Ask not why the top-quark is so heavy—ask why the other quarks are so light."

Indeed, in fermionic models one has a scheme to try to explain the hierarchy of fermion masses, as non-zero Yukawa couplings are rather restricted by the many stringy selection rules at play [32, 15]. Typically at the cubic level only some fermions acquire Yukawa couplings (*e.g.*, those of the third generation). Yukawa couplings for the lighter generations appear at higher-order in non-renormalizable interactions,

$$\begin{array}{ll} \lambda Q_3 t^c H & \lambda_t \sim g \\ \lambda Q_2 c^c H \frac{\langle\phi\rangle}{M} & \lambda_c \sim g \frac{\langle\phi\rangle}{M} \\ \lambda Q_1 u^c H \frac{\langle\phi\rangle^2}{M^2} & \lambda_u \sim g \frac{\langle\phi\rangle^2}{M^2} \end{array} \tag{46}$$

A hierarchy of effective Yukawa couplings is generated once the ratios $\langle\phi\rangle/M \sim 1/10$ (from $\mathrm{U_A}(1)$ cancellation) are inserted. In practice this exercise is complicated by the question of how to embed the Standard Model fields in the string representations. Moreover, the resulting Yukawa matrix need not be diagonal. Taking all these effects into account it is possible to obtain semi-realistic fermion mass spectra [15, 20].

By far the simplest Yukawa coupling to calculate is that of the top quark

$$\lambda_t(M_{Pl}) = (g\sqrt{2})\left(\frac{g}{\sqrt{2}}\right) = g^2 \approx 0.7 \tag{47}$$

where we have inserted a normalization factor $(g/\sqrt{2})$ coming from the Kähler function normalization [23], and the value of g from running the Standard Model gauge couplings up to the string scale. The top-quark mass itself,

$$m_t = \left(\frac{v}{\sqrt{2}}\right)\lambda_t(m_t)\frac{\tan\beta}{\sqrt{\tan^2\beta+1}}\ , \tag{48}$$

involves dynamics: running from $M_{Pl} \to M_Z$ to obtain $\lambda_t(M_Z)$, and at M_Z the ratio of Higgs vevs ($\tan\beta$). (There is also a +7% QCD correction to the running mass to obtain the experimentally observable "pole" mass.) The result for m_t as a function of $\tan\beta$ is shown in Fig. 5, from where we conclude that

$$m_t \approx (160 - 190)\,\mathrm{GeV}\ , \tag{49}$$

a result in good agreement with experimental observations. Such first-principles predictions should not be confused with similar numerical results obtained in SO(10) GUT models, where the values of the b and τ masses are used to deduce m_t.

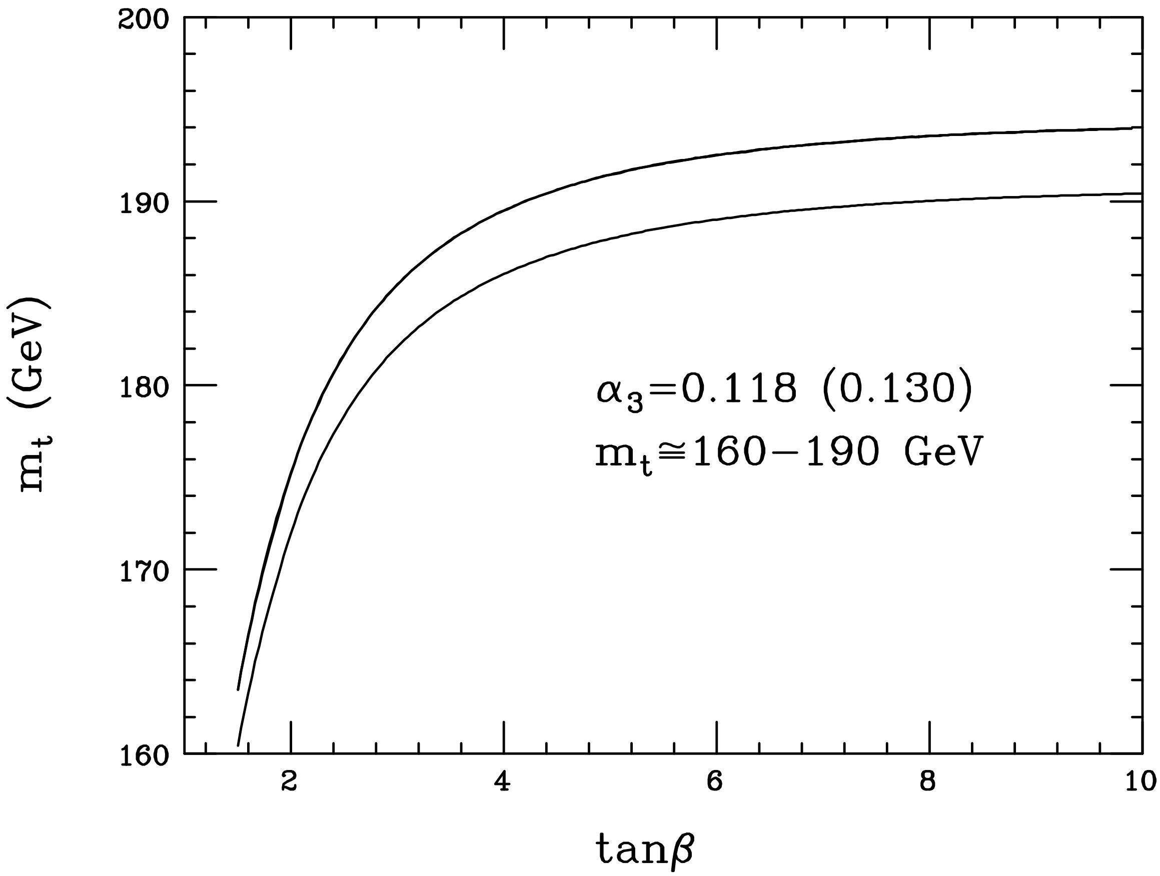

Figure 5: The calculated value of the top-quark mass as a function of $\tan\beta$ for $\alpha_3(M_Z) = 0.118\,(0.130)$ (lower (upper) curve).

2.4 String unification: a new scenario

Typical string models possess gauge groups which are products of various factors. In fact, all of the corresponding gauge couplings are predicted to be equal [33] at the string scale [34]

$$M_{\rm string} = 5 \times g \times 10^{17}\,{\rm GeV}\ . \tag{50}$$

(This statement holds for gauge groups realized with level-one Kac-Moody algebras, as in the case in our model of interest.) The crucial point is that the defining property of GUTs, namely the unification of the gauge couplings at and above the GUT scale, is obtained in string models without appealing to traditional GUT structures. These models have then been called Grand Unified Superstring Theories or GUSTs [19].

In the case of flipped SU(5), two scenarios for string unification may be considered, depending on where the SU(5)×U(1) gauge symmetry breaks down to the Standard Model gauge group. In either scenario, the mechanism for generating the GUT-symmetry breaking vevs is the vacuum shift that results in the cancellation of the anomalous $U_A(1)$. In the original scenario [35], one assumed that these vevs were as large as possible consistent with this cancellation mechanism, with the GUT

symmetry breaking occuring at the string scale. This scenario appears quite reasonable as the string unification scale (50) nearly coincides with the scale from $\mathrm{U_A}(1)$ cancellation (39), which is assumed to provide the SU(5)×U(1)-breaking vevs. Recent work considers a slightly different scenario [36], wherein SU(5)×U(1) breaks near the traditional GUT scale (M_{32}) and SU(5) and U(1) eventually "superunify" at the string scale. This scale is more non-trivial to derive, and our mechanism plays a more unique role. Such range of vevs are perfectly allowed, and essentially determine the flat direction along which SU(5)×U(1) breaks.[3] The new scenario may also be more appealing to some, as $M_{32} \ll M_{\rm string}$, making effective field theory calculations more reliable.

The recent "two-step" scenario has some interesting consequences not present in the original "one-step" scenario. First of all, in either scenario one must consider string models with an extra pair of ($\mathbf{10},\overline{\mathbf{10}}$) SU(5) representations to possibly achieve string unification. (This statement has been recently re-examined in Ref. [38], and shown to be unavoidable.) In the two-step scenario the ($\mathbf{10},\overline{\mathbf{10}}$) mass scale ($M_{10}$) needs to be [36]

$$M_{10} \sim 10^{8-9}\,\mathrm{GeV}\ . \tag{51}$$

On the other hand, in the one-step scenario one must split the $Q, \bar{Q}$ from the $D^c, \bar{D}^c$ pieces in the $\mathbf{10},\overline{\mathbf{10}}$, giving them masses of order 10^{12} GeV and 10^{6} GeV respectively [35]. In either scenario these scales are in principle calculable from the scales of hidden matter condensation, although in the two-step scenario one needs to generate only one such scale. The breaking of SU(5)×U(1) is assumed to occur in both scenarios in the shifting vacuum process; in the two-step scenario this requires a smaller vev, which is certainly possible (see first row in Table 4). The two-step scenario also makes use of the apparent "LEP" scale ($M_{\rm LEP} \approx 10^{16}$ GeV), since SU(3) and SU(2) unify at this scale irrespective of the presence of the extra ($\mathbf{10},\overline{\mathbf{10}}$), as complete SU(5) multiplets do not affect this result.[4]

The running of the Standard Model gauge couplings in both two- and one-step scenarios are shown in Fig. 6. The dotted lines in the one-step scenario correspond to the case of minimal SU(5) unification. Note that to lowest order, in the two-step scenario, $M_{32} \sim M_{\rm LEP}$ and the SU(3) beta function vanishes above the M_{10} scale. Also, in the actual string model one expects the appearance of fractionally charged particles in the hidden sector which are likely to modify the running of the gauge couplings. For our present purposes we have assumed that all these particles acquire masses of the order $\sqrt{\epsilon} \approx M_{\rm string}$, and thus do not contribute to the gauge coupling RGEs. In a full analysis such effects will have to be consistently included.

[3]It is interesting to point out that in attempts at constructing traditional GUT models in strings (via higher-level Kac-Moody algebras) [37], the generation of the GUT scale has remained unclear. If such models ever prove to be realistic, and contain an anomalous $\mathrm{U_A}(1)$ factor, then our mechanism should offer a possibility to obtain the GUT scale.

[4]Note that if one wants to obtain the apparent LEP scale in a string model, SU(3) and SU(2) need to be unified at this scale, a result not possible in either standard-like SU(3)×SU(2)×U(1) [39, 20], or Pati-Salam SU(4)×SU(2)×SU(2) [30], or $\mathrm{SU}(3)^3$ [40] type string models, as these string-unify SU(3) and SU(2) at $M_{\rm string} \gg M_{\rm LEP}$.

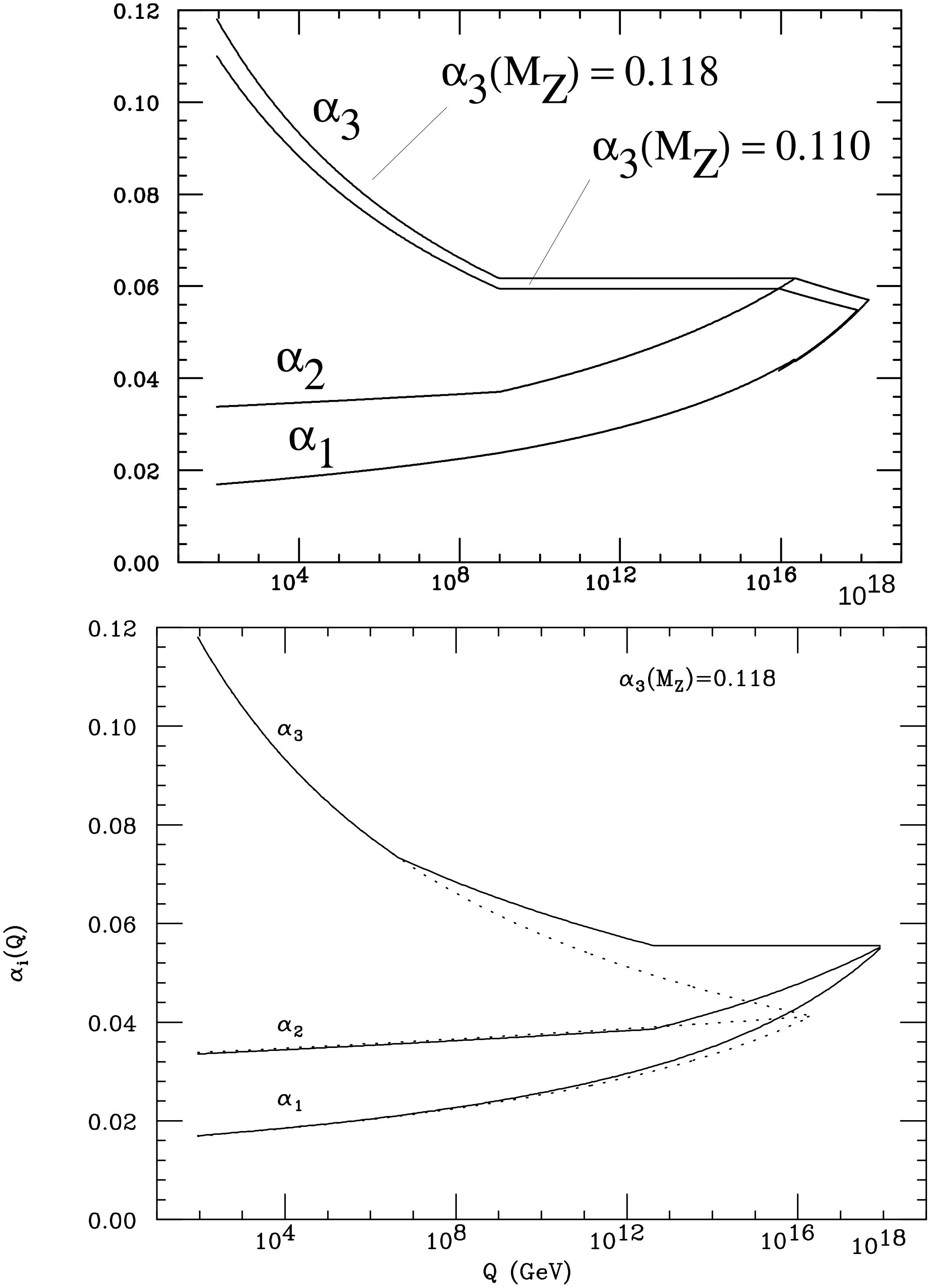

Figure 6: Two-step and one-step scenarios for string unification in flipped SU(5).

3 Experimental Flipped SU(5)

As we have seen above, string models are quite predictive, with for instance all superparticle masses ($\propto m_{3/2}$) and fermion masses calculable. However, one should realize that predictions do vary from model to model. Moreover, string models are complicated: one must choose one out of many-many possible models, they have hidden and observable sectors, they have unwanted fields that must decouple, the various intermediate scales need to be generated dynamically, etc. Nonetheless, models can be worked out and predictions can be derived. In this section we would like to describe a possible scenario for a flipped SU(5) string model, as pertains to the superparticle masses. We warn the reader that other scenarios may exist, depending on how one assigns the Standard Model particles to the string representations. All of these possible scenarios are expected to yield predictions for the top-quark mass within the range in Eq. (49), although the required value of $\tan\beta$ and the precise m_t prediction may differ from scenario to scenario.

3.1 A possible scenario [40]

In this scenario one obtains for the supersymmetry-breaking masses

- Gaugino masses (universal)

$$m_{1/2} = m_{3/2} \tag{52}$$

- Scalar masses

$$\begin{aligned}
\text{First generation} \quad & m^2_{Q_1,U^c_1,D^c_1,L_1,E^c_1} = 0 && (53)\\
\text{Second generation} \quad & m^2_{Q_2,U^c_2,D^c_2,L_2,E^c_2} = 0 && (54)\\
\text{Third generation} \quad & m^2_{Q_3,D^c_3} = m^2_{3/2},\ m^2_{U^c_3,L_3,E^c_3} = 0 && (55)\\
\text{Higgs masses} \quad & m^2_{H_1} = m^2_{H_2} = 0 && (56)
\end{aligned}$$

- Trilinear scalar couplings (universal)

$$A_0 = m_{3/2} \tag{57}$$

- Bilinear scalar coupling

$$B_0 = m_{3/2} \tag{58}$$

Our scenario has only one free parameter. At high energies our unknowns are $m_{3/2}$ and μ.[5] At low-energies we inherit the ratio of Higgs vacuum expectation values ($\tan\beta$), which must exceed unity for the radiative electroweak breaking mechanism to work. This mechanism imposes two constraints on the model parameters, which

[5] In principle, even μ is calculable in this class of models. We expect to obtain an effective μ term from non-renormalizable contributions to the superpotential, *i.e.*, $\lambda h\bar{h}\langle TT\rangle/M \Rightarrow \mu = \lambda\langle TT\rangle/M$.

can be used to determine $|\mu|$ and $\tan\beta$ in terms of $m_{3/2}$. Thus, our model can be described in terms of a single parameter: $m_{1/2} \leftrightarrow m_{\tilde{g}} \leftrightarrow m_{\chi_1^\pm}$

To obtain the low-energy spectrum, starting from the high-energy supersymmetry-breaking mass parameters in Eqs. (52-58), we need to follow the standard procedure of running the coupled renormalization group equations for all masses, gauge, and Yukawa couplings. The following analysis has been carried out in the one-step scenario for string unification, as discussed in Sec. 2.4. Calculations for the new two-step scenario are underway [36]. At the electroweak scale we minimize the one-loop effective potential and determine $|\mu|$ and B (see *e.g.*, Ref. [42]). We adjust $\tan\beta$ until the computed B_0 agrees with first-principles calculation of B_0 in Eq. (58). One finds that only

$$\mu < 0 \tag{59}$$

works, and that the values of $\tan\beta$, even though $m_{3/2}$ dependent, fall in the following narrow range

$$\tan\beta \approx 2.2 - 2.3\ . \tag{60}$$

Therefore, from Eq. (48) one can determine m_t

$$m_t \approx 175\,(179)\,\mathrm{GeV}, \quad \alpha_3 = 0.118\,(0.130)\ . \tag{61}$$

This result is shown in Fig. 7.

The determination of the full spectrum of superparticle masses then follows, although not without subtlety. Because the non-universality of scalar masses at the string scale, the renormalization group equations receive a new contribution, which effectively amounts to a shift in all scalar mass parameters [43]

$$\Delta m_i^2 = -c^2 Y_i f\ , \tag{62}$$

where Y_i is the hypercharge of the particle, $f = 0.0060$ is an RGE factor, and c^2 is the non-universality coefficient at the string scale

$$c^2 = m_{H_2}^2 - m_{H_1}^2 + \sum_{i=1,2,3}\left(m_{Q_i}^2 + m_{D_i^c}^2 + m_{E_i^c}^2 - m_{L_i}^2 - 2m_{U_i^c}^2\right)\ . \tag{63}$$

In our case $c^2 = 2m_{1/2}^2$. This shift affects all sparticle masses, but is significant only for the right-handed sleptons, as these possess only hypercharge quantum numbers. The right-handed slepton masses are given by

$$m_{\tilde{\ell}_R}^2 = a\,m_{1/2}^2 + \tan^2\theta_W M_W^2 \frac{\tan^2\beta - 1}{\tan^2\beta + 1} \tag{64}$$

with $a = 0.153$ in the usual universal case, but $a = 0.153 - 0.120 = 0.034$ here. This mass should be compared with that of the lightest neutralino χ_1^0

$$m_{\chi_1^0} \approx 0.25 m_{1/2} \tag{65}$$

$$m_{\tilde{\ell}_R} \approx \sqrt{(0.18 m_{1/2})^2 + (36)^2}\ , \tag{66}$$

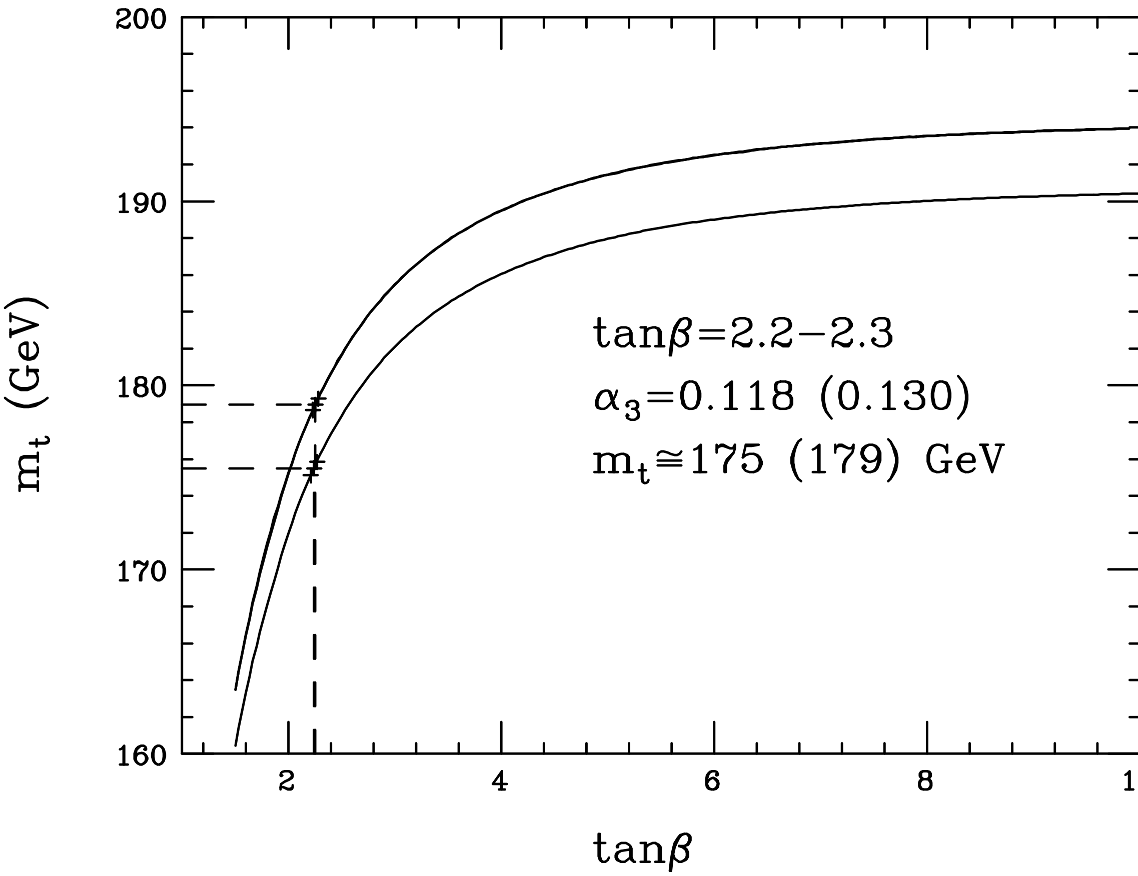

Figure 7: The calculated value of the top-quark mass as a function of $\tan\beta$ for $\alpha_3(M_Z) = 0.118\,(0.130)$ (lower (upper) curve), indicating the predicted value of $\tan\beta$, and therefore of the top-quark mass.

which shows that if $m_{1/2}$ is too large, then $\tilde{\ell}_R$ becomes the lightest supersymmetric particle (LSP). Since such electrically charged particle would be stable, our model cannot support such regime, and we obtain an important cutoff in the one-dimensional parameter space

$$m_{1/2} \lesssim 180\,\mathrm{GeV}\ . \tag{67}$$

The full spectrum is shown in Fig. 8. Of note are the following features

$$m_{\chi_1^\pm} \approx m_{\chi_2^0} < 90\,\mathrm{GeV} \tag{68}$$

$$m_h < 90\,\mathrm{GeV} \tag{69}$$

$$m_{\tilde{\ell}_R} < 50\,\mathrm{GeV} \tag{70}$$

$$m_{\tilde{q}} \approx 0.98 m_{\tilde{g}} \tag{71}$$

$$\tilde{t}_{1,2}\ \text{big split} \tag{72}$$

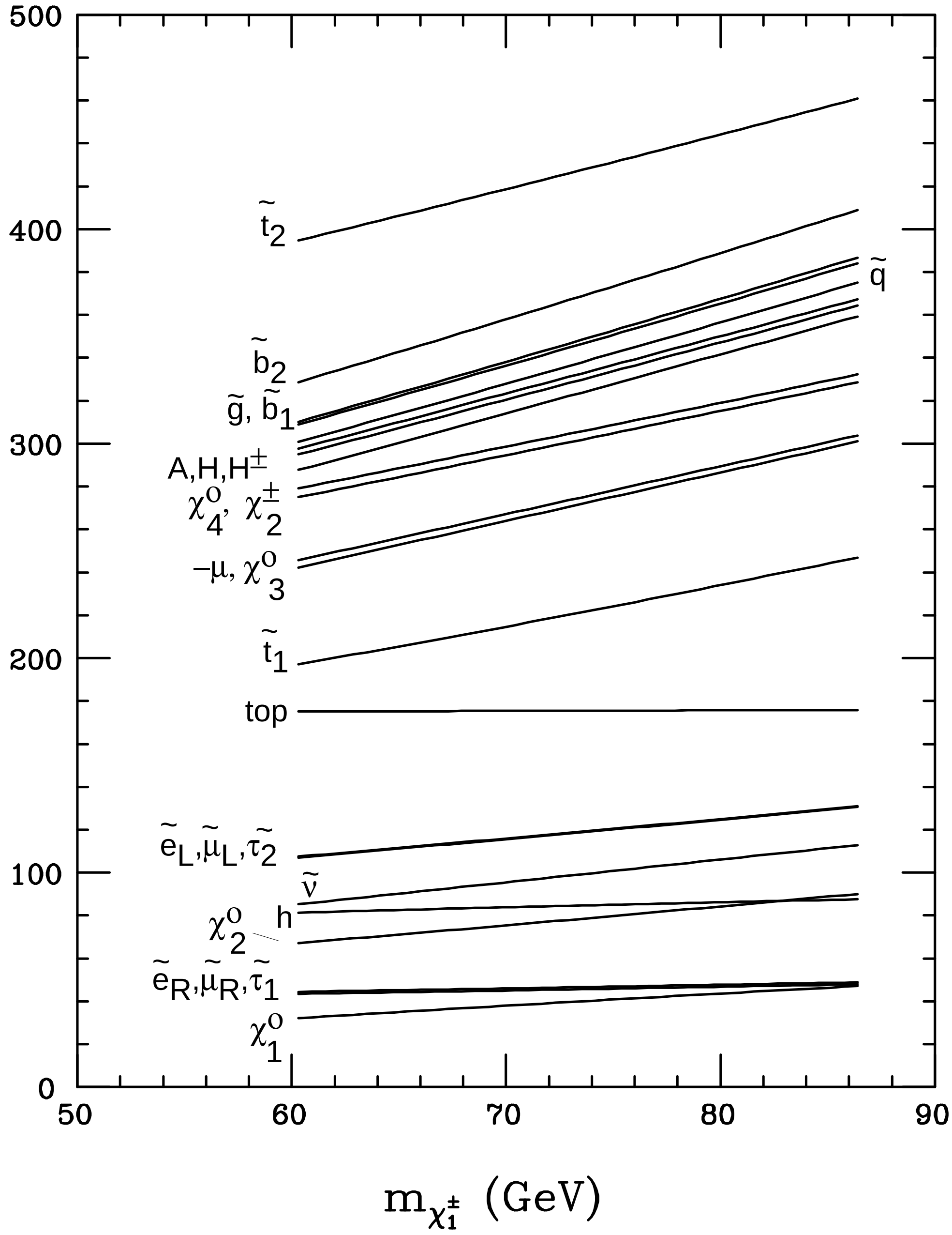

Figure 8: Full spectrum of superparticles in a possible scenario of stringy flipped SU(5). Note the cutoff in the parameter space resulting from demanding a neutral and colorless lightest supersymmetric particle (LSP).

3.2 Discovery prospects at the Tevatron

The spectrum of the model, as shown in Fig. 8 indicates that the traditional supersymmetry signals at hadron colliders, namely searches for squarks and gluinos via the missing-energy channel, are not accessible at the present-day Tevatron, as one has $m_{\tilde{q},\tilde{g}} \approx (300-400)\,\mathrm{GeV}$. Some of this parameter space may become accessible in the Main Injector era, and more so in the case of a high-luminosity upgrade of the Tevatron [44]. This mass range would be, however, easily accessible at the LHC.

A much more promising and immediate avenue for discovery is afforded by considering the production and decay of neutralinos and charginos, as the lightest of these must satisfy $m_{\chi^0_2,\chi^\pm_1} < 90\,\mathrm{GeV}$. The reactions of interest are [45, 46]

$$p\bar{p} \to \chi^0_2\chi^\pm_1 X \to 3\ell \quad \text{(trileptons)} \tag{73}$$
$$p\bar{p} \to \chi^+_1\chi^-_1 X \to 2\ell \quad \text{(dileptons)} \tag{74}$$

The corresponding branching ratios can be calculated, as everything else, in terms of our one parameter, giving [41]

$$B(\chi^\pm_1 \to \ell^\pm) \approx 1/2 \quad (\ell = e+\mu) \tag{75}$$
$$B(\chi^\pm_1 \to 2j) \approx 1/4 \tag{76}$$
$$B(\chi^0_2 \to \ell^+\ell^-) = 2/3 \tag{77}$$

(Note that the last ratio is maximal, as this channel proceeds via the two-body decay $\chi^0_2 \to \tilde{\ell}_R\ell$). These sizeable branching ratios and relatively light superparticles imply significant trilepton and dilepton rates. These are shown in Fig. 9, along with an estimate of the experimental sensitivity expected by the end of Run IB ($\int\mathcal{L} = 100\,\mathrm{pb}^{-1}$) [46]. Clearly, complete exploration of this model at the Tevatron via trilepton events is guaranteed.

Searches for the light sleptons predicted in this model ($m_{\tilde{\ell}_R} < 50\,\mathrm{GeV}$) appear also possible via $p\bar{p} \to \tilde{\ell}^+_R\tilde{\ell}^-_R X$. This process has a cross section of $1\,(0.1)\,\mathrm{pb}$ for $m_{\tilde{\ell}_R} = 45\,(50)\,\mathrm{GeV}$, as we expect $B(\tilde{\ell}_R \to \ell\chi^0_1) = 1$ into dileptons. However, the prospects for slepton detection are not very bright, given the low rates and large backgrounds that are expected.

3.3 Discovery prospects at LEPII

At LEPII, starting to operate in the Fall of 1995, one could detect the lightest Higgs boson, the lightest chargino, and the right-handed sleptons. Discovery of the latter two should be straightforward, with only a slight increase in the beam energy and a modest amount of luminosity.

Charginos are predicted to obey $m_{\chi^\pm_1} \approx (60-90)\,\mathrm{GeV}$, and should be most easily detectable via [47]

$$e^+e^- \to \chi^+_1\chi^-_1 \to \ell + 2j + \text{missing} \tag{78}$$

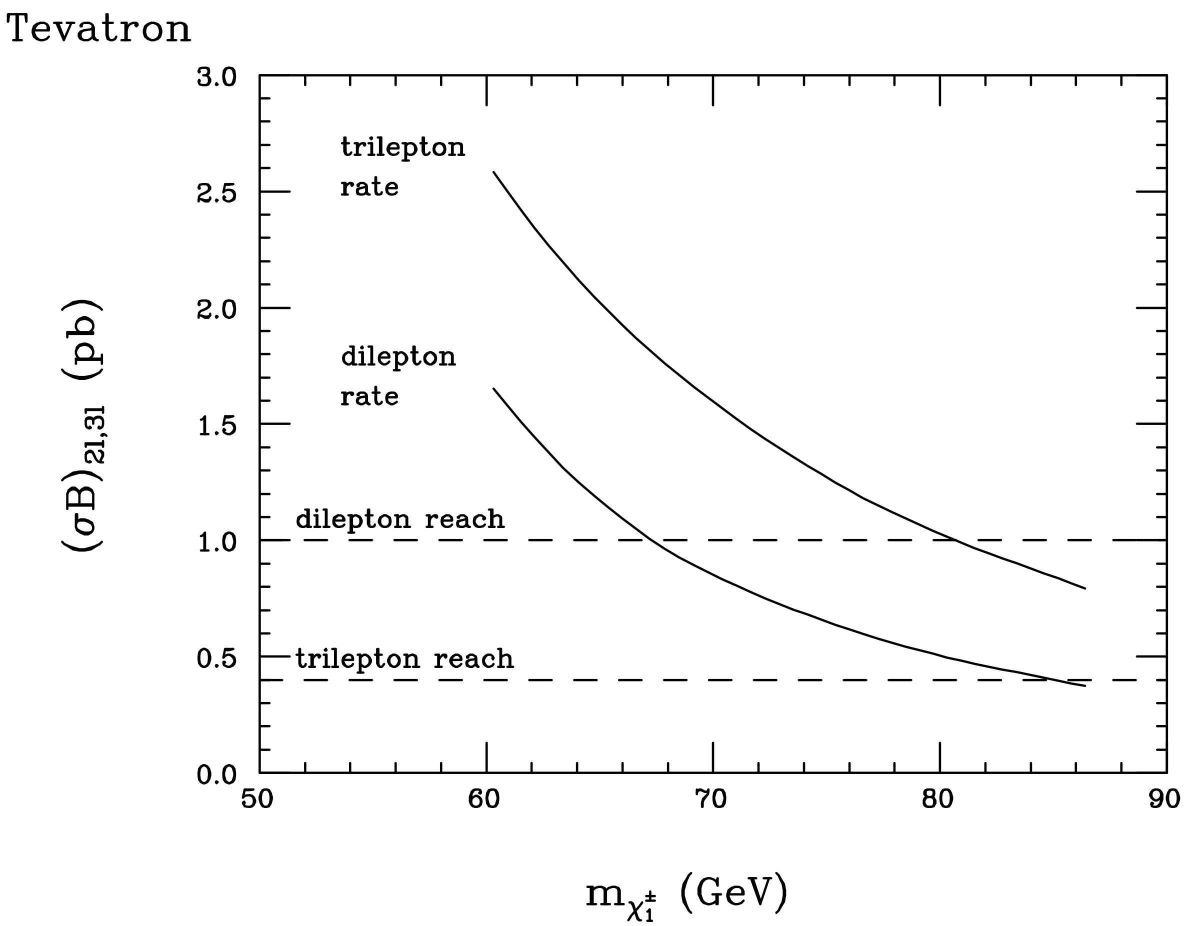

Figure 9: Calculated trilepton and dilepton rates from chargino-neutralino production at the Tevatron versus the chargino mass. Also indicated are the expected experimental sensitivities with $100\,\mathrm{pb}^{-1}$ of accumulated data.

Inserting the chargino branching ratios given above, for $m_{\chi_1^\pm} = 60\,(90)$ GeV we expect a cross section of $2.1\,(0.62)$ pb (at $\sqrt{s} = 200\,\mathrm{GeV}$). With a sufficiently high beam energy, chargino production into the "mixed" mode should be quite noticeable. Right-handed sleptons must be rather light in this model ($m_{\tilde{\ell}_R} < 50\,\mathrm{GeV}$) and will be pair-produced

$$e^+e^- \to \tilde{\ell}_R^+ \tilde{\ell}_R^- \to 2\ell + \mathrm{missing} \tag{79}$$

We expect $\sigma(\tilde{e}_R) > 2\,\mathrm{pb}$ and $\sigma(\tilde{\mu}_R) > 0.9\,\mathrm{pb}$ (at $\sqrt{s} = 200\,\mathrm{GeV}$), which should make detection immediate. Finally, because of the prediction $\tan\beta \approx 2.2$, the Higgs boson will be at the low end of its possible range ($m_h \approx 80 - 90\,\mathrm{GeV}$), and should be produced via [48]

$$e^+e^- \to Z^* \to Zh \to (f\bar{f})(b\bar{b}) \tag{80}$$

Such Higgs boson will be indistinguishable from the Standard Model Higgs boson ($\sin^2(\alpha - \beta) = 0.996 - 0.998$). The estimated mass reach of LEPII as a function of

the center-of-mass energy [49],

$$m_h \approx \sqrt{s} - 95\,\mathrm{GeV}\ , \tag{81}$$

will make it sensitive to our Higgs boson for $\sqrt{s} \gtrsim (175 - 185)\,\mathrm{GeV}$. This is not expected to occur in the initial phase of the LEPII upgrade. It is worth remarking that the supersymmetry channel $h \to \chi^0_1\chi^0_1$ erodes the usual $h \to b\bar{b}$ signal a little.

3.4 Discovery prospects via Rare processes

We consider three rare processes through which this model could be tested indirectly: $b \to s\gamma$ at CLEO, $(\mathrm{g}-2)_\mu$ at Brookhaven, and R_b at LEP.

3.4.1 $\mathrm{B}(b \to s\gamma)$

This decay mode has been observed by the CLEO Collaboration with the following result [50]

$$\mathrm{B}(b \to s\gamma)^{\mathrm{exp}} = (1-4) \times 10^{-4}\ . \tag{82}$$

This result agrees well with the Standard Model prediction (for all allowed values of m_t) and therefore constrains any extension of the Standard Model. In particular, in supersymmetry $\mathrm{B}(b \to s\gamma)$ varies a lot over the multi-dimensional parameter space [51]. In fact, it is possible to find regions of parameter space where the supersymmetry prediction is larger or smaller than the Standard Model predicition by more than two orders of magnitude. The dominant diagram involves the chargino–top-squark loop and depends strongly on $\tan\beta$. Also, QCD corrections are large, and to date only fully known to leading order. In the present model we find the following range [41]

$$\mathrm{B}(b \to s\gamma) = [(4.2-5.3) \to (3.9-5.1)] \times 10^{-4}\ , \tag{83}$$

which is in fair agreement with the experimental result in Eq. (82). (The uncertainties due to uncalculated next-to-leading order QCD corrections have been estimated by varying the renormalization scale around the b-quark mass [52].) It should be noted that our small value of $\tan\beta$ helps in suppressing the supersymmetric contributions sufficiently, given the lightness of our spectrum.

3.4.2 $(\mathrm{g}-2)_\mu$

The anomalous magnetic moment of the muon was last measured in 1970, with the result $a_\mu^{\mathrm{exp}} = 1165923\,(8.5)\times 10^{-9}$ [53]. When this result is contrasted with the present Standard Model prediction ($a_\mu^{\mathrm{SM}} = 1165919.20\,(1.76) \times 10^{-9}$ [54]), it allows a 95%CL interval for new physics contributions [55]

$$-13.2 \times 10^{-9} < a_\mu^{susy} < 20.8 \times 10^{-9}\ . \tag{84}$$

There are two supersymmetric contributions to a_μ, with charginos and sneutrinos in the loop, or with sleptons and neutralinos in the loop. The latter is small because of

the small $\tilde{\mu}_L - \tilde{\mu}_R$ mixing angle ($\propto m_\mu$). The dominant chargino-sneutrino contribution is greatly enhanced by large values of $\tan\beta$ (not unlike the chargino–top-squark contribution in $b \to s\gamma$), and can lead to values of $a_\mu^{\rm susy}$ outside the allowed range in Eq. (84) [55]. Nonetheless we find [41]

$$a_\mu^{\rm susy} = (-2.4 \to -1.7) \times 10^{-9} \tag{85}$$

which is within the present limits. The new E821 experiment at Brookhaven, scheduled to start taking date in 1996, should achieve a sensitivity of 0.4×10^{-9} [56] and thus be greatly sensitive to this prediction.

3.4.3 R_b

The ratio $R_b = \Gamma(Z \to b\bar{b})/\Gamma(Z \to \text{hadrons})$ has been measured at LEP to be $R_b^{\rm exp} = 0.2219 \pm 0.0017$ [57], which is more than three standard deviations above the Standard Model prediction of $R_b^{\rm SM} = 0.2157$. We obtain [41]

$$R_b^{\rm susy} = (4.4 \to 3.2) \times 10^{-4} \ , \tag{86}$$

which shifts the Standard Model prediction in the direction of the experimental result only slightly. This shift will not be observable unless the experimental sensitivity increases by a factor of four, and then only if the experimental result has somehow been reconciled with the Standard Model prediction.

3.5 Prospects for Dark Matter detection

In our model the lightest supersymmetric particle (LSP) is stable (as R-parity is unbroken) and will contribute to the cold dark matter in the Universe. The calculation of the relic abundance of neutralinos yields [41]

$$\Omega_\chi h^2 \lesssim 0.025 \ , \tag{87}$$

where h is the scaled Hubble parameter. Such magnitude of neutralino relic density is of interest in models with a sizeable cosmological constant [58], where one could expect

$$\Omega_\nu(0.3) + \Omega_\chi(0.1) + \Omega_\Lambda(0.6) = 1 \ , \tag{88}$$

with a hot dark matter component (neutrinos, see Sec. 1.2.5), a cold dark matter component (neutralinos), and a cosmological constant component. Our neutralinos would populate the galactic halo and could be detected via scattering off nuclei [59]. In present-day cryogenic Germanium detectors, the present sentivity is of 0.1 events/kg/day (eventually expected to improve to 0.01 events/kg/day). In Fig. 10 we show the relic density of neutralinos as a function of the LSP mass, showing a dip when the higgs-boson pole is encountered. We also show the predicted detection rate in the Ge detector, with its characteristic kinematical peak when the LSP mass is half of the Ge mass.

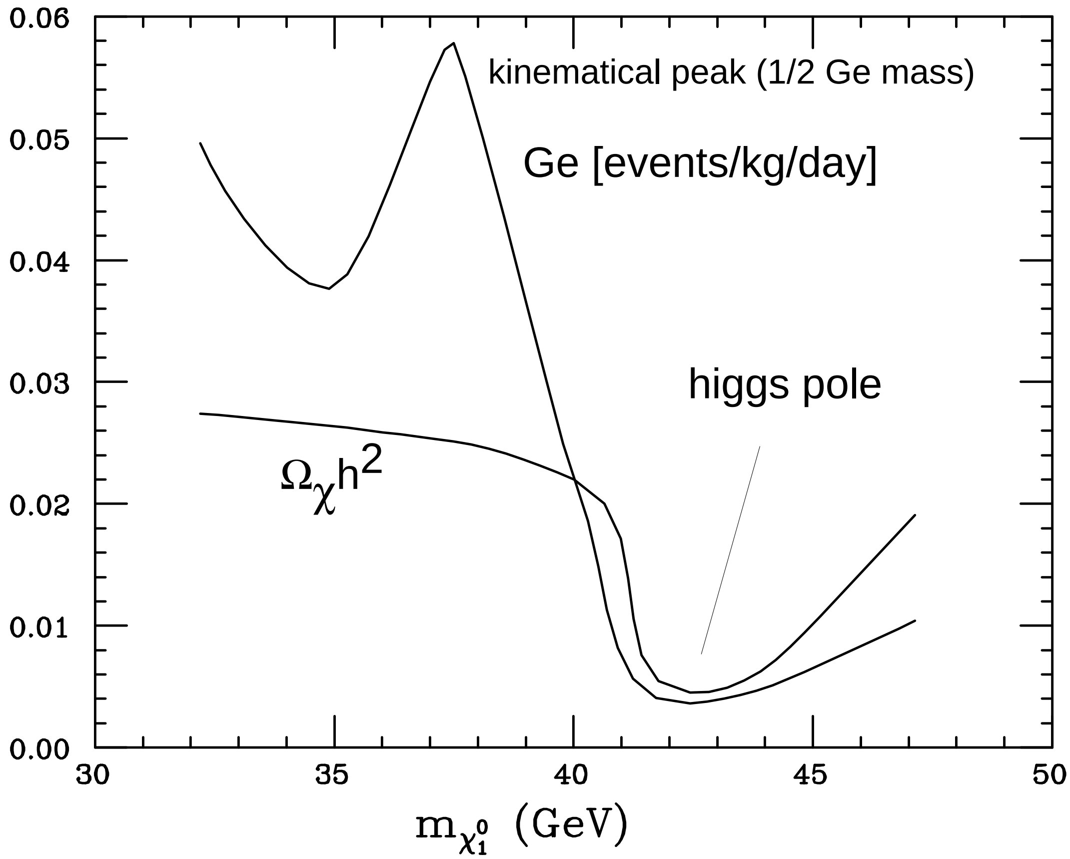

Figure 10: The relic density of neutralinos as a function of the LSP mass. Also shown is the predicted detection rate in the Ge detector in events/kg/day.

4 Conclusions

We have reviewed the features that make flipped SU(5) and very nice supersymmetric unified theory, including recent developments that allow a natural reduction in the prediction for α_s, which is consistent with the whole range of presently experimentally allowed values. The real power of flipped SU(5) lies, however, in its being string derivable, as then many more predictions follow. This string model suppresses naturally the tree-level and one-loop contributions to the vacuum energy, and has a distinct prediction for the supersymmetry-breaking parameters. The anomalous $\mathrm{U_A}(1)$ plays a crucial role in the one-loop suppression, and is conjectured to suppress the vacuum energy to all orders in string perturbation theory. Another new development consists of a new scenario for flipped string unification, wherein the "LEP" scale (where the SU(3) and SU(2) gauge couplings unify) is obtained by means of having the SU(5)×U(1) breaking vevs participate in the anomalous $\mathrm{U_A}(1)$ cancellation mechanism. The SU(5) and U(1) gauge couplings then unify at the string scale. This

scenario predicts the existence of a new $(\mathbf{10},\overline{\mathbf{10}})$ complete pair at the scale 10^{8-9} GeV, that can be obtained from hidden sector matter condensation. Among the tests of flipped SU(5) we have its prediction for α_s and the proton lifetime into the $e^+\pi^0$ mode, its prediction for the top-quark mass, its prediction for the superparticle masses and the rates at which they could be produced in direct and indirect processes, and its prediction for the hot (ν_τ) and cold (χ^0_1) dark matter in the Universe and its detection.

Acknowledgments

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